

# Light Stop Decay in the MSSM with Minimal Flavour Violation

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## Abstract

In supersymmetric scenarios with a light stop particle  $\tilde{t}_1$  and a small mass difference to the lightest supersymmetric particle (LSP) assumed to be the lightest neutralino, the flavour changing neutral current decay  $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$  can be the dominant decay channel and can exceed the four-body stop decay for certain parameter values. In the framework of Minimal Flavour Violation (MFV) this decay is CKM-suppressed, thus inducing long stop lifetimes. Stop decay length measurements at the LHC can then be exploited to test models with minimal flavour breaking through Standard Model Yukawa couplings. The decay width has been given some time ago by an approximate formula, which takes into account the leading logarithms of the MFV scale. In this paper we calculate the exact one-loop decay width in the framework of MFV. The comparison with the approximate result exhibits deviations of the order of 10% for large MFV scales due to the neglected non-logarithmic terms in the approximate decay formula. The difference in the branching ratios is negligible. The large logarithms have to be resummed. The resummation is performed by the solution of the renormalization group equations. The comparison of the exact one-loop result and the tree level flavour changing neutral current decay, which incorporates the resummed logarithms, demonstrates that the resummation effects are important and should be taken into account.

# 1 Introduction

The Standard Model (SM) provides a very successful effective theory of particle interactions which is in excellent agreement with electroweak precision data. Furthermore, remarkable consistency and precision tests have been made in the sector of quark flavour violation. These tests and limits on flavour changing neutral currents (FCNC) from  $K$ ,  $D$  and  $B$  studies put strong constraints on possible New Physics beyond the SM [1]. They forbid a generic flavour structure of New Physics at the TeV scale and raise the question why contributions from New Physics at  $\sim 1$  TeV are strongly suppressed. A solution to this New Physics Flavour Puzzle is provided by the framework of Minimal Flavour Violation (MFV) [2–5]. It requires all sources of flavour and CP-violation to be given by the SM structure of the Yukawa couplings. Flavour mixing in models of New Physics is then always proportional to the off-diagonal elements of the Cabibbo-Kobayashi-Maskawa (CKM) matrix [6]. In particular, the mixing of the third generation squarks with the first and second generation squarks is highly suppressed by small CKM quark mixing angles.

Since the flavour structure of New Physics at the TeV scale must be non-generic, flavour measurements provide a good probe of New Physics. One of the best studied examples is given by supersymmetry (SUSY). In the context of flavour violation the phenomenology of the light stop partner  $\tilde{t}_1$  is especially interesting. In most SUSY models a light stop quark mass arises naturally. Due to the large top Yukawa coupling the mixing between the weak eigenstates  $\tilde{t}_L$  and  $\tilde{t}_R$  leads to a large mass splitting between the stop mass eigenstates. Furthermore, the large top Yukawa coupling in general drives the stop mass via the renormalization group equation (RGE) running to smaller masses, even if all squarks have a common mass at the SUSY breaking scale. In scenarios with a light stop which predominantly decays into a charm quark and the lightest neutralino, assumed to be the lightest supersymmetric particle (LSP),  $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$ , squark flavour violation can be tested by exploiting stop decay length measurements [7]. It has been shown, that in these scenarios light stops can be discovered at the LHC [8]. Finally, a light stop is also favoured by baryogenesis which requires a stop mass with about the top mass value or less for successful electroweak baryogenesis within the context of the Minimal Supersymmetric Extension of the Standard Model (MSSM) [9].

Minimal Flavour Violation naturally arises in supergravity models which provide flavour-independent scalar mass terms at a high scale like the Planck scale  $M_P$ . Some time ago the authors of Ref. [10] have provided an approximate formula for the flavour changing neutral current (FCNC) decay  $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$  by starting from a vanishing tree level  $\tilde{t}_1 - c - \tilde{\chi}_1^0$  coupling at the Planck scale. The decay then proceeds via one-loop diagrams. The non-vanishing divergences, which are due to scalar self-mass diagrams, have been subtracted by a soft counterterm at the Planck scale so that a large logarithm  $\ln(M_P^2/M_W^2)$  remains at the weak scale chosen to be  $M_W$ . The authors argue that in view of this large logarithm, the remaining non-logarithmic part of the one-loop diagrams can hence safely be neglected so that their result for the decay width takes a rather simple form.

In this work, we perform the complete one-loop calculation of the  $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$  decay in the framework of MFV. We perform the full renormalization program and keep the finite non-logarithmic terms arising from the loop integrals. This allows us to study the importance of the neglected non-logarithmic pieces in the formula given by Ref. [10].

To get a reliable result, the appearing large logarithms should be resummed. In the renor-

malization group approach this corresponds to solving the renormalization group equations for the scalar soft SUSY breaking squark masses. As has been pointed out in [4], the hypothesis of MFV is not renormalization group invariant. Flavour off-diagonal squark mass terms are hence induced by the Yukawa couplings, so that the squark and quark mass matrices cannot be diagonalized simultaneously any more and the stop state receives some admixture from the charm squark, inducing a FCNC between stop, charm and LSP neutralino,  $\tilde{t}_1 - c - \tilde{\chi}_1^0$ . From this point of view, the logarithmic piece of our one-loop result is equivalent to the first order in the expansion of the RGE solution for the squark-quark-neutralino coupling in powers of  $\alpha$ , whereas the tree level decay calculated with the FCNC coupling includes the resummation of the large logarithms. The comparison of the two decay widths provides an estimate of the importance of the resummation of the large logarithms.

The outline of our paper is as follows: In section 2 we present the diagrams contributing to the one-loop decay. We set up our notation for the squark and quark sector in section 3. The counterterms and the renormalization are discussed in section 4. Section 5 contains the numerical analysis. We conclude in section 6. In the Appendix we list our Feynman rules, the various amplitudes contributing to the decay and we derive the FCNC counterterm.

## 2 One-loop decay

We work in the framework of the MSSM with MFV so that all flavour changing effects with quarks and squarks are controlled by the quark Yukawa couplings and CKM mixing angles [3]. The decay of the lightest stop  $\tilde{t}_1$  into the lightest neutralino  $\tilde{\chi}_1^0$  and a charm quark  $c$ ,

$$\tilde{t}_1 \rightarrow c \tilde{\chi}_1^0, \quad (1)$$

is then mediated at the one-loop level. We consider scenarios where the light stop  $\tilde{t}_1$  is the next-to-lightest supersymmetric particle (NLSP) and the lightest neutralino is the LSP. The process is built up by the stop and charm self-energies and the vertex diagrams, *cf.* Fig. 1. Note, that in our calculation we set

$$m_c = 0. \quad (2)$$

Therefore in the  $\tilde{t}_1$  self-energies we have only non-vanishing contributions for transitions into the left-handed charm squark  $\tilde{c}_L$ . Those into right-handed scharm,  $\tilde{c}_R$ , are zero for  $m_c = 0$ . All diagrams are mediated by charged current loops. The various diagrams which contribute are depicted in Fig. 2. The self-energies and vertex corrections are divergent and have to be renormalized. The counterterms for the squark and quark self-energies and for the vertex renormalization are shown in Fig. 3. The FCNC vertex does not arise at tree level. Its occurrence as counterterm at one-loop

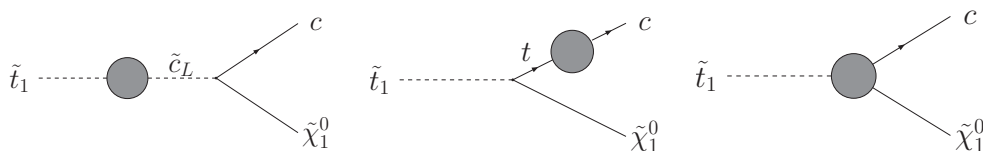


Figure 1: Generic diagrams contributing to the loop-decay  $\tilde{t}_1 \rightarrow c \tilde{\chi}_1^0$ .

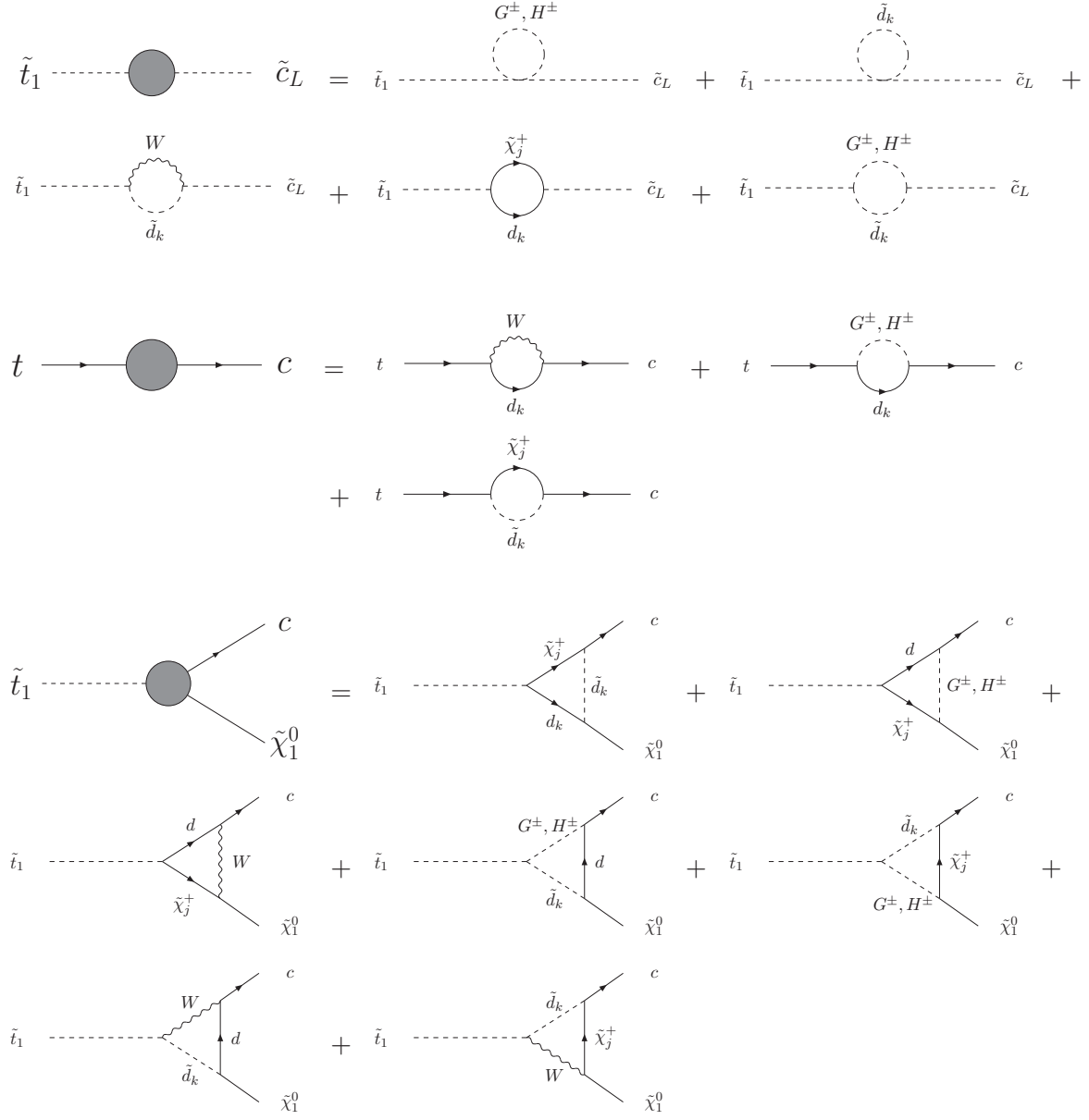


Figure 2: Generic diagrams contributing to the squark and quark self-energy and the proper vertex correction.

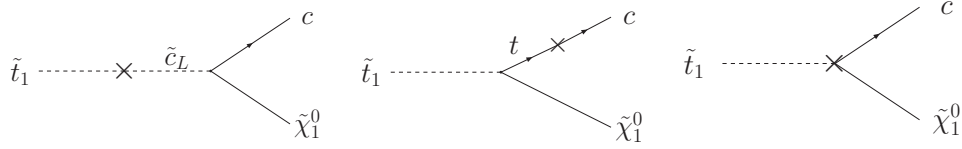


Figure 3: Counterterm diagrams.

level is due to the fact that MFV is not RGE-invariant, since the weak interactions affect the squark and quark mass matrices differently [11]. Their simultaneous diagonalization cannot be maintained at higher orders so that it can consistently be imposed at a single scale only, called  $\mu_{\text{MFV}}$  in the following.

For the calculation of the stop decay process we define an effective interaction vertex

$$T \equiv g \bar{u}_c(k_2) (F_L \mathcal{P}_L + F_R \mathcal{P}_R) v_{\tilde{\chi}_1^0}(k_1) , \quad (3)$$

where  $\bar{u}_c, v_{\tilde{\chi}_1^0}$  denote the charm and neutralino spinors and  $k_1, k_2$  are the four-momenta of the outgoing neutralino and charm quark.  $F_L$  and  $F_R$  are form factors associated with the chirality projectors  $\mathcal{P}_L$  and  $\mathcal{P}_R$ , respectively. They receive contributions  $F^v$  from the vertex diagrams,  $F^{tc}, F^{\tilde{t}_1 \tilde{c}_L}$  from the quark and squark self-energies, and  $\delta F^v, \delta F^{tc}, \delta F^{\tilde{t}_1 \tilde{c}_L}$  from the vertex and the quark and squark wave function renormalization counterterms,

$$F_{L,R} = [F^{\tilde{t}_1 \tilde{c}_L} + \delta F^{\tilde{t}_1 \tilde{c}_L} + F^{tc} + \delta F^{tc} + F^v + \delta F^v]_{L,R} . \quad (4)$$

They are specified in section 4 and the Appendices B1-3 and C.

### 3 The quark and squark sector

For the choice of our notation the quark and squark sectors are discussed here in more detail. The definition of the couplings is deferred to Appendix A. We define the  $3 \times 3$  unitary matrices  $U^{u_{L,R}}, U^{d_{L,R}}$  as the matrices which rotate the left- and right-handed up- and down-type quark current eigenstates  $u_{L,R}, d_{L,R}$  to their corresponding mass eigenstates,  $u_{L,R}^m, d_{L,R}^m$ ,

$$u_L^m = U^{u_L} u_L, \quad u_R^m = U^{u_R} u_R, \quad d_L^m = U^{d_L} d_L, \quad d_R^m = U^{d_R} d_R . \quad (5)$$

The CKM matrix  $\mathcal{V}$  is given by

$$\mathcal{V} = U^{u_L} U^{d_L \dagger} . \quad (6)$$

For the squark interaction eigenstate we define a six component vector

$$\tilde{q}' = \begin{pmatrix} \tilde{q}'_L \\ \tilde{q}'_R \end{pmatrix} , \quad (7)$$

where  $\tilde{q}'_L, \tilde{q}'_R$  are each a three component column vector in generation space. The squared squark mass matrix can be written as a  $2 \times 2$  Hermitian matrix of  $3 \times 3$  blocks

$$\mathcal{M}_{\tilde{q}'}^2 = \begin{pmatrix} \mathcal{M}_{\tilde{q}'_{LL}}^2 & \mathcal{M}_{\tilde{q}'_{LR}}^2 \\ \mathcal{M}_{\tilde{q}'_{RL}}^2 & \mathcal{M}_{\tilde{q}'_{RR}}^2 \end{pmatrix} . \quad (8)$$

It is diagonalized by a  $6 \times 6$  unitary matrix  $\tilde{W}^q$  which rotates the squark interaction eigenstates to their mass eigenstates  $\tilde{q}^m$ ,

$$\tilde{q}^m = \tilde{W}^q \tilde{q}' . \quad (9)$$

The six component column vector  $\tilde{q}^m$  is defined to be ordered in mass, with  $\tilde{q}_1$  being the lightest squark. Equation (9) can then be rewritten as

$$\begin{aligned}\tilde{q}_s^m &= \tilde{W}_{si} \tilde{q}'_{iL} + \tilde{W}_{s\,i+3} \tilde{q}'_{iR} \quad (s = 1, \dots, 6, \, i = 1, 2, 3) \\ &\equiv (\tilde{W}_L \tilde{q}'_L + \tilde{W}_R \tilde{q}'_R)_s ,\end{aligned}\tag{10}$$

where  $i$  denotes the generation index.<sup>1</sup> The rotation of the squarks by the same unitary matrices  $U^{qL,R}$  as the quarks defines the super-CKM basis. In models with non-minimal flavour violation the squark mass matrix is flavour-mixed in this basis, in contrast to the quark mass matrix. In models with MFV at the scale  $\mu_{\text{MFV}}$ , however, the squarks can be rotated by  $U^{qL,R}$  to their flavour eigenstates in parallel to the quarks, and the super-CKM basis is at the same time the flavour eigenstate basis. Hence, suppressing generation indices,

$$\tilde{q}_L = U^{qL} \tilde{q}'_L, \quad \tilde{q}_R = U^{qR} \tilde{q}'_R .\tag{11}$$

The squared mass matrix in the flavour eigenstate basis  $(\tilde{q}_L, \tilde{q}_R)^T$  then reads

$$\mathcal{M}_{\tilde{q}}^2 = \begin{pmatrix} (\tilde{M}_{\tilde{q}_L}^2 + m_q^2) \mathbf{1}_3 & m_q (A_q - \mu r_q) \mathbf{1}_3 \\ m_q (A_q - \mu r_q) \mathbf{1}_3 & (\tilde{M}_{\tilde{q}_R}^2 + m_q^2) \mathbf{1}_3 \end{pmatrix} ,\tag{12}$$

where  $r_d = 1/r_u = \tan \beta$  for down- and up-type squarks. With  $\tan \beta$  we denote the ratio of the vacuum expectation values of the two complex Higgs doublets. They are introduced in order to generate masses of up- and down-type fermions [12]. The parameter  $A_q$  denotes the trilinear coupling of the soft SUSY breaking part of the Lagrangian,  $\mu$  the Higgsino mass parameter and  $m_q$  the quark partner mass.  $\mathbf{1}_3$  is a  $3 \times 3$  unit matrix in generation space. The parameters  $\tilde{M}_{\tilde{q}_{L,R}}$  are given by the left- and right-handed scalar soft SUSY breaking masses  $M_{\tilde{q}_{L,R}}$  and the  $D$ -terms,

$$\begin{aligned}\tilde{M}_{\tilde{q}_{L,R}}^2 &= M_{\tilde{q}_{L,R}}^2 + D_{\tilde{q}_{L,R}} \\ D_{\tilde{q}_L} &= M_Z^2 \cos 2\beta (I_q^3 - Q_q \sin^2 \theta_W) \\ D_{\tilde{q}_R} &= M_Z^2 \cos 2\beta Q_q \sin^2 \theta_W ,\end{aligned}\tag{13}$$

where  $I_q^3$  denotes the third component of the weak isospin,  $Q_q$  the electric charge,  $M_Z$  the  $Z$  boson mass and  $\theta_W$  the Weinberg angle. The squared mass matrix  $\mathcal{M}_{\tilde{q}}^2$  can be diagonalized by a  $6 \times 6$  unitary matrix  $W^q$  which rotates the flavour eigenstates to their mass eigenstates,

$$\begin{aligned}\tilde{q}_s^m &= W_{st} \begin{pmatrix} \tilde{q}_L \\ \tilde{q}_R \end{pmatrix}_t = W_{si} \tilde{q}_{Li} + W_{s\,i+3} \tilde{q}_{Ri} \equiv (W_L \tilde{q}_L + W_R \tilde{q}_R)_s \\ &\quad (s, t = 1, \dots, 6, \, i = 1, 2, 3) .\end{aligned}\tag{14}$$

Comparison with Eqs. (10) and (11) shows that we can factorize the  $6 \times 3$  matrices  $\tilde{W}_{L,R}$  into the  $6 \times 3$  flavour-diagonal matrices  $W_{L,R}$  and the  $3 \times 3$  quark rotation matrices defined above,

$$\tilde{W}_L = W_L U^{qL} \quad \text{and} \quad \tilde{W}_R = W_R U^{qR} ,\tag{15}$$

with  $q = u, d$ . The matrix  $W$  can be expressed in terms of mixing angles by

$$(W_L)_{ii} = (W_R)_{i+3\,i} = \cos \theta_{q_i} \quad , \quad (W_R)_{ii} = -(W_L)_{i+3\,i} = \sin \theta_{q_i} .\tag{16}$$

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<sup>1</sup>We have thus decomposed the mass eigenstate squark field into left- and right-chiral interaction eigenstate squarks.

For the three quark generations  $i$  the relation between the flavour eigenstates  $\tilde{q}_{iL}, \tilde{q}_{iR}$  and the squark mass eigenstates  $\tilde{q}_s^m = (\tilde{q}_i^m, \tilde{q}_{i+3}^m)$  hence reads

$$\begin{aligned}\tilde{q}_i^m &= \tilde{q}_{iL} \cos \theta_{q_i} + \tilde{q}_{iR} \sin \theta_{q_i} \\ \tilde{q}_{i+3}^m &= -\tilde{q}_{iL} \sin \theta_{q_i} + \tilde{q}_{iR} \cos \theta_{q_i} .\end{aligned}\quad (17)$$

For better legibility, we suppress the generation indices from now on wherever possible, and the lighter and heavier squark mass eigenstates are generically called  $\tilde{q}_1$  and  $\tilde{q}_2$ . The mixing angles are then given by

$$\sin 2\theta_q = \frac{2m_q(A_q - \mu r_q)}{M_{\tilde{q}_1}^2 - M_{\tilde{q}_2}^2} \quad , \quad \cos 2\theta_q = \frac{\tilde{M}_{\tilde{q}_L}^2 - \tilde{M}_{\tilde{q}_R}^2}{M_{\tilde{q}_1}^2 - M_{\tilde{q}_2}^2} , \quad (18)$$

and the masses of the squark mass eigenstates read

$$M_{\tilde{q}_{1,2}}^2 = m_q^2 + \frac{1}{2} \left[ \tilde{M}_{\tilde{q}_L}^2 + \tilde{M}_{\tilde{q}_R}^2 \mp \sqrt{(\tilde{M}_{\tilde{q}_L}^2 - \tilde{M}_{\tilde{q}_R}^2)^2 + 4m_q^2(A_q - \mu r_q)^2} \right] . \quad (19)$$

Since the mixing angles are proportional to the masses of the quarks, the mixing is important in the stop sector and can drive the lightest stop mass even lighter than the top quark mass.

## 4 Counterterms and Renormalization

The quark and squark self-energies and the vertex diagrams are ultraviolet (UV) divergent and need to be renormalized. The quark and squark wave functions are renormalized on-shell. With the bare (s)quark fields  $q^{(0)}$  ( $\tilde{q}^{(0)}$ ) related to the renormalized (s)quark fields  $q$  ( $\tilde{q}$ ) by

$$\tilde{q}^{(0)} = \left(1 + \frac{1}{2}\delta Z^{\tilde{q}}\right) \tilde{q} \quad \text{and} \quad q_{L,R}^{(0)} = \left(1 + \frac{1}{2}\delta Z^{L,R}\right) q_{L,R} , \quad (20)$$

this leads for the squarks to the following off-diagonal elements of the wave function renormalization constants in terms of the real part of the squark self-energy  $\tilde{\Sigma}$ ,

$$\delta Z_{st}^{\tilde{q}} = \frac{2}{m_{\tilde{q}_s}^2 - m_{\tilde{q}_t}^2} \text{Re} \tilde{\Sigma}_{st}(m_{\tilde{q}_t}^2) \quad s, t = 1, \dots, 6, \quad s \neq t . \quad (21)$$

Defining the following structure for the quark self-energy,

$$\Sigma_{ij}(p^2) \equiv \not{p} \Sigma_{ij}^L(p^2) \mathcal{P}_L + \not{p} \Sigma_{ij}^R(p^2) \mathcal{P}_R + m_i \Sigma_{ij}^{Ls}(p^2) \mathcal{P}_L + m_j \Sigma_{ij}^{Rs}(p^2) \mathcal{P}_R , \quad (22)$$

we have for the corrections to the off-diagonal chiral components of the quark wave functions,

$$\begin{aligned}\delta Z_{ij}^L &= \frac{2}{m_{q_i}^2 - m_{q_j}^2} \left[ m_{q_i}^2 \text{Re} \Sigma_{ij}^{Ls}(m_{q_j}^2) + m_{q_j}^2 \text{Re} \Sigma_{ij}^{Rs}(m_{q_j}^2) + m_{q_j}^2 \text{Re} \Sigma_{ij}^L(m_{q_j}^2) + m_{q_i} m_{q_j} \text{Re} \Sigma_{ij}^R(m_{q_j}^2) \right] \\ \delta Z_{ij}^R &= \frac{2}{m_{q_i}^2 - m_{q_j}^2} \left[ m_{q_i} m_{q_j} \text{Re} \Sigma_{ij}^{Ls}(m_{q_j}^2) + m_{q_i} m_{q_j} \text{Re} \Sigma_{ij}^{Rs}(m_{q_j}^2) + m_{q_i} m_{q_j} \text{Re} \Sigma_{ij}^L(m_{q_j}^2) + \right. \\ &\quad \left. m_{q_j}^2 \text{Re} \Sigma_{ij}^R(m_{q_j}^2) \right] \quad i, j = 1, 2, 3, \quad i \neq j .\end{aligned}\quad (23)$$

As for our scenario we chose the  $\tilde{t}_1$  to be the NLSP, all our self-energies are real, and after the on-shell renormalization of the quark and squark wave functions we are only left with the one-loop vertex diagrams and the FCNC vertex counterterm. From now on “Re” will be dropped. We

regularize the divergences by dimensional regularization in  $n = 4 - 2\epsilon$  dimensions so that ultraviolet singularities appear as poles in  $\epsilon$ . We have explicitly verified that the same result is obtained with dimensional reduction. By exploiting the unitarity relations of the CKM matrix as well as those of the chargino mixing matrices  $U$  and  $V$ , defined in Appendix A, we find for the vertex contribution to the form factors

$$g F_R^v = -ig \mathcal{F} \left[ \frac{1}{\epsilon} + \ln \frac{\bar{\mu}^2}{m_{\text{loop}}^2} + \text{finite terms} \right] \quad (24)$$

$$g F_L^v = 0, \quad (25)$$

where  $\bar{\mu}$  denotes the 't Hooft mass of dimensional regularization. We have used the short-hand notation

$$\mathcal{F} \equiv \frac{1}{16\pi^2} g^2 \sqrt{2} \left[ \frac{Z_{11}}{6} t_W + \frac{Z_{12}}{2} \right] \left( \frac{\mathcal{V}_{cb} \mathcal{V}_{tb}^* m_b^2 \cos \theta_t}{2M_W^2 c_\beta^2} \right), \quad (26)$$

where  $Z_{11}, Z_{12}$  denote matrix elements of the  $4 \times 4$   $Z$  matrix, which diagonalizes the neutralino mass matrix, *cf.* Appendix A. The further 'finite terms' in Eq. (24), which do not depend on  $\ln \bar{\mu}^2$ , can be extracted from the full one-loop result for the vertex which is given in Appendix B3. We have introduced a generic mass for the loop particles,  $m_{\text{loop}}$ . Note, that in the numerical analysis we will use the exact results with the different loop particle masses. Here, for reasons of legibility and also to make later contact with the result derived in Ref. [10] we adopt the generic notation. The left-handed form factor  $F_L$  is zero due to our choice of vanishing  $c$ -quark mass.

The FCNC counterterm arises from the flavour non-diagonal part of the wave function renormalization, from the renormalization of the quark and squark mixing matrices [13–16] and the renormalization of the quark masses.<sup>2</sup> The renormalized squark mixing matrix  $\tilde{W}^r$  is related to the bare  $\tilde{W}^{(0)}$  by

$$\tilde{W}_{su}^{(0)} = (\delta_{st} + \delta \tilde{w}_{st}) \tilde{W}_{tu}^r, \quad s, t, u = 1, \dots, 6. \quad (27)$$

And similarly, the renormalized quark mixing matrices  $U^{uL,Rr}$  are related to the bare quark matrices  $U^{uL,R(0)}$  by

$$U_{ik}^{uL,R(0)} = (\delta_{ij} + \delta u_{ij}^{uL,R}) U_{jk}^{uL,Rr}, \quad i, j, k = 1, 2, 3. \quad (28)$$

The indices  $s, t, u$  denote the six squark mass eigenstates, and  $i, j, k$  are generation indices. We impose the MFV condition on the renormalized mixing matrices  $\tilde{W}^r, U^r$  and hence demand them to be flavour diagonal. This leads to flavour non-diagonal counterterms  $\delta \tilde{w}, \delta u^{uL,R}$ . Furthermore, as the bare and renormalized mixing matrices are unitary the counterterms must be antihermitian. The UV divergent part of each counterterm is determined such that it cancels the divergent part of the antihermitian part of the corresponding wave function renormalization matrix [14–16],

$$\delta \tilde{w} = \frac{1}{4} (\delta Z^{\tilde{q}} - \delta Z^{\tilde{q}\dagger}) \quad (29)$$

$$\delta u^{uL,R} = \frac{1}{4} (\delta Z^{L,R} - \delta Z^{L,R\dagger}). \quad (30)$$

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<sup>2</sup>In our renormalization procedure and the definition of the counterterm we follow the same approach as in Ref. [17].



The general form of the FCNC vertex counterterm depicted in Fig. 3 has been derived in Appendix C. For our process we have the following form factor contributions to the vertex counterterm,

$$g \delta F_R^v = -ig e_{L1}^u \left[ \frac{1}{2} \delta Z_{ct}^{L\dagger} \cos \theta_t + \frac{1}{2} \delta Z_{\tilde{c}_L \tilde{t}_1}^{\tilde{q}} + \delta u_{ct}^{U_L} \cos \theta_t + \delta \tilde{w}_{\tilde{c}_L \tilde{t}_1}^\dagger \right] \quad (31)$$

$$g \delta F_L^v = 0, \quad (32)$$

with

$$e_{L1}^u = \sqrt{2} \left[ \frac{Z_{11}}{6} \tan \theta_W + \frac{1}{2} Z_{12} \right] \quad (33)$$

and

$$\delta Z_{ct}^{L\dagger} = 2\Sigma_{tc}^{Ls}(m_c^2 = 0) = 0 \quad (34)$$

$$\delta Z_{\tilde{c}_L \tilde{t}_1}^{\tilde{q}} = \frac{2\Sigma_{\tilde{c}_L \tilde{t}_1}^{\tilde{q}}(m_{\tilde{t}_1}^2)}{m_{\tilde{c}_L}^2 - m_{\tilde{t}_1}^2}. \quad (35)$$

The squark wave function has been renormalized at  $p^2 = m_{\tilde{t}_1}^2$ . The finite parts of the counterterms Eqs. (29,30) depend on the renormalization conditions. Absorbing also the finite part of the antihermitian wave function renormalization leads to a gauge-dependent on-shell renormalization scheme [15, 16, 18]. Performing minimal subtraction on the other hand is gauge-independent [19] and imposes the MFV condition on the  $\overline{\text{MS}}$  parameters at the scale  $\mu_{\text{MFV}}$ . In the following we will adopt this scheme. Consequently, the result will depend on the MFV scale  $\mu_{\text{MFV}}$ . The squark mixing matrix counterterm then reads<sup>3</sup>

$$\delta \tilde{w}_{\tilde{c}_L \tilde{t}_1}^\dagger = \frac{1}{2} \left( \frac{\Sigma_{\tilde{c}_L \tilde{t}_1}^{\tilde{q}}(m_{\tilde{c}_L}^2) + \Sigma_{\tilde{c}_L \tilde{t}_1}^{\tilde{q}}(m_{\tilde{t}_1}^2)}{m_{\tilde{t}_1}^2 - m_{\tilde{c}_L}^2} \right)_{\overline{\text{MS}}}. \quad (36)$$

A gauge invariant prescription for the quark mixing matrix counterterm is given by [15],

$$\delta u_{ct}^{U_L} = -\frac{1}{2} [\Sigma_{tc}^L(0) + 2\Sigma_{tc}^{Ls}(0)]_{\overline{\text{MS}}} = -\frac{1}{2} \Sigma_{tc}^L(0)_{\overline{\text{MS}}}. \quad (37)$$

For the quark mixing matrix contribution to the vertex counterterm we then find

$$-g e_{L1}^u \cos \theta_t \left( \delta u_{ct}^{U_L} \right)_{\overline{\text{MS}}} = g \frac{\mathcal{F}}{2} \left[ \frac{1}{\epsilon} + \ln \frac{\bar{\mu}^2}{\mu_{\text{MFV}}^2} \right]. \quad (38)$$

For the contribution from the squark mixing we have

$$-g e_{L1}^u \left( \delta \tilde{w}_{\tilde{c}_L \tilde{t}_1}^\dagger \right)_{\overline{\text{MS}}} = g \frac{\mathcal{F}}{2} \left( \frac{-m_{\tilde{t}_1}^2 - m_{\tilde{c}_L}^2 - 2\mathcal{A}}{m_{\tilde{t}_1}^2 - m_{\tilde{c}_L}^2} \right) \left[ \frac{1}{\epsilon} + \ln \frac{\bar{\mu}^2}{\mu_{\text{MFV}}^2} \right], \quad (39)$$

with

$$\mathcal{A} = -\mu^2 + A_b^2 + \tilde{M}_{bR}^2 + c_\beta^2 (M_W^2 (t_\beta^2 - 1) + M_A^2 t_\beta^2) + m_t A_b \tan \theta_t. \quad (40)$$

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<sup>3</sup>This counterterm definition can lead to large contributions to the matrix element if  $m_{\tilde{t}_1} \approx m_{\tilde{c}_L}$ . For a discussion, see *e.g.* Refs. [20]. In the scenarios of our numerical analysis we have  $m_{\tilde{t}_1} \ll m_{\tilde{c}_L}$ .

It depends on the Higgsino parameter  $\mu$ , the soft SUSY breaking mass parameter  $\tilde{M}_{\tilde{b}_R}$  including  $D$  term contributions, the trilinear coupling  $A_b$ , the mixing angle  $\beta$ , the  $W$  boson mass and the pseudoscalar Higgs mass  $M_A$ .

The contribution from  $\delta Z_{\tilde{c}_L \tilde{t}_1}^{\tilde{q}}$  is given by

$$-g e_{L1}^u \frac{1}{2} \left( \delta Z_{\tilde{c}_L \tilde{t}_1}^{\tilde{q}} \right) = g \frac{\mathcal{F}}{2} \left( \frac{2m_{\tilde{t}_1}^2 + 2\mathcal{A}}{m_{\tilde{t}_1}^2 - m_{\tilde{c}_L}^2} \right) \left[ \frac{1}{\epsilon} + \ln \frac{\bar{\mu}^2}{m_{\text{loop}}^2} + \text{finite terms} \right]. \quad (41)$$

As before, we have introduced a generic loop particle mass, and 'finite terms' denote further terms which do not depend on  $\ln \bar{\mu}^2$ . Their specific form can be extracted from the explicit formulae of the stop self-energies given in Appendix B1. Inserting Eqs. (34,38,39,41) in Eq. (31), the right-chiral part of the FCNC counterterm is then given by

$$g \delta F_R^v = ig \mathcal{F} \left[ \frac{1}{\epsilon} - \frac{m_{\tilde{c}_L}^2 + \mathcal{A}}{m_{\tilde{t}_1}^2 - m_{\tilde{c}_L}^2} \ln \frac{\bar{\mu}^2}{\mu_{\text{MFV}}^2} + \frac{m_{\tilde{t}_1}^2 + \mathcal{A}}{m_{\tilde{t}_1}^2 - m_{\tilde{c}_L}^2} \ln \frac{\bar{\mu}^2}{m_{\text{loop}}^2} + \text{finite terms} \right]. \quad (42)$$

Adding Eqs. (24) and (42) and replacing  $\mathcal{F}$  and  $\mathcal{A}$  by Eq. (26) and Eq. (40), respectively, we arrive at the following final result for the form factors, which contribute to Eq. (4),

$$g F_R = \frac{i}{16\pi^2} g^3 \sqrt{2} \left[ \frac{Z_{11}}{6} t_W + \frac{Z_{12}}{2} \right] \left( \frac{\mathcal{V}_{cb} \mathcal{V}_{tb}^* m_b^2 \cos \theta_t}{2M_W^2 c_\beta^2} \right) \left( \frac{1}{m_{\tilde{t}_1}^2 - m_{\tilde{c}_L}^2} \right) \times \\ [m_{\tilde{c}_L}^2 - \mu^2 + A_b^2 + \tilde{M}_{\tilde{b}_R}^2 + c_\beta^2 (M_W^2 (t_\beta^2 - 1) + M_A^2 t_\beta^2) + m_t A_b \tan \theta_t] \ln \left( \frac{\mu_{\text{MFV}}^2}{m_{\text{loop}}^2} \right) \\ + \text{finite terms} \quad (43)$$

$$g F_L = 0. \quad (44)$$

Finally, the stop decay width in terms of the form factors Eqs. (43,44) is given by

$$\Gamma(\tilde{t}_1 \rightarrow c \tilde{\chi}_1^0) = \frac{g^2 m_{\tilde{t}_1}}{16\pi} \left( 1 - \frac{m_{\tilde{\chi}_1^0}^2}{m_{\tilde{t}_1}^2} \right)^2 |F_R|^2. \quad (45)$$

As can be inferred from  $F_R$ , depending on the scale of MFV, the logarithm can become very large, and the decay can become important in certain regions of the parameter space, especially for large values of  $\tan \beta$ . The finite terms, which do not depend on  $\ln \mu_{\text{MFV}}^2$ , are then only subleading. If we drop the finite terms in Eq. (43), the approximate result given by Hikasa and Kobayashi in Ref. [10] should be reproduced. In fact, for  $m_c = 0$  we can rewrite

$$m_{\tilde{c}_L}^2 - \mu^2 + c_\beta^2 (M_W^2 (t_\beta^2 - 1) + M_A^2 t_\beta^2) = M_{H_d}^2 + M_{\tilde{q}_L}^2 + \frac{1}{3} M_Z^2 \sin^2 \theta_W \cos 2\beta, \quad (46)$$

where  $M_{H_d}$  denotes the mass parameter of the Higgs doublet  $H_d$  which couples to down-type fermions. With this relation the form factor Eq. (43) leads to the approximate result  $F_R^{H/K}$  of Ref. [10], if we set the MFV scale equal to the Planck scale,  $\mu_{\text{MFV}} = M_P$ , choose  $M_W$  as generic loop particle mass and neglect the finite terms,

$$g F_R^{H/K} = \frac{i}{16\pi^2} g^3 \sqrt{2} \left[ \frac{Z_{11}}{6} t_W + \frac{Z_{12}}{2} \right] \left( \frac{\mathcal{V}_{cb} \mathcal{V}_{tb}^* m_b^2 \cos \theta_t}{2M_W^2 c_\beta^2} \right) \left( \frac{1}{m_{\tilde{t}_1}^2 - m_{\tilde{c}_L}^2} \right) \times \\ [M_{H_d}^2 + M_{\tilde{q}_L}^2 + A_b^2 + M_{\tilde{b}_R}^2 + m_t A_b \tan \theta_t] \ln \left( \frac{M_P^2}{M_W^2} \right). \quad (47)$$

In our full one-loop calculation of the decay width in terms of  $F_R$ , Eq. (43), the finite terms are included, and the relevance of these contributions can be checked by comparing with the approximate result  $\Gamma^{H/K}$  for the decay width in terms of the form factor  $F_R^{H/K}$ , Eq. (47). This will be discussed in section 5.

To get a reliable result, the large logarithms of the MFV scale in the decay formula should be resummed. The logarithm is related to the running of the FCNC coupling of the neutralino to a quark and squark of different generations. We have required this coupling to vanish at the scale  $\mu_{\text{MFV}}$ . Minimal Flavour Violation is not RGE-invariant, however. Even though MFV is imposed at  $\mu_{\text{MFV}}$ , at any other scale  $\mu \neq \mu_{\text{MFV}}$  a FCNC coupling will be generated through renormalization group evolution. The solution of the one-loop RGE for the quark and squark mixing matrices provides the resummation of the large  $\ln \mu_{\text{MFV}}^2$ . The coefficient of this logarithm in Eq. (43) is then given by the first order in the expansion of the RGE solution for the squark-quark-neutralino coupling in powers of  $\alpha$ . In the following we will call the right-handed form factor including the resummation effects  $F_R^{FV}$ . It is given by the FCNC coupling obtained through renormalization group evolution including flavour violation of the squark and quark mixing matrices, from some high scale down to the scale relevant for the decay process.

## 5 Numerical Analysis

The scenarios for the numerical analysis have been chosen such that they lead to a NLSP stop  $\tilde{t}_1$  and a  $\tilde{\chi}_1^0$  LSP. The latter represents a promising dark matter (DM) candidate in the MSSM [21]. The mass difference is chosen to be small enough so that the loop mediated flavour changing decay  $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$  is dominating and can compete with the four-body decays<sup>4</sup> into the LSP, a  $b$ -quark and a fermion pair [23],

$$\tilde{t}_1 \rightarrow \tilde{\chi}_1^0 b f \bar{f}' . \quad (48)$$

Such scenarios can be consistent with electroweak baryogenesis [9] and also with Dark Matter constraints [24]. The mass spectra and mixing angles have been calculated with the spectrum calculator SPheno [25] and compared to SOFTSUSY [26]. Both codes include the option to perform two-loop RGE running with and without the inclusion of flavour violation and both support the SUSY Les Houches accord [27]. Within this accord the gauge and Yukawa couplings as well as the soft SUSY breaking mass parameters and trilinear couplings are given out as  $\overline{\text{DR}}$  running parameters at a scale  $Q$ , which we have chosen to be the scale of electroweak symmetry breaking (EWSB). We have verified that the calculation of the decay width leads to the same result if we apply dimensional reduction instead of dimensional regularization, so that the  $\overline{\text{DR}}$  running parameters can be used. The mixing matrix elements and the SUSY particle pole masses have been taken at the scale of EWSB as well. The SM parameters have been chosen as  $M_Z = 91.187$  GeV,  $\alpha_{em}^{-1\overline{\text{MS}}}(M_Z) = 127.934$ ,  $\alpha_s^{\overline{\text{MS}}}(M_Z) = 0.1184$ ,  $m_b^{\overline{\text{MS}}}(m_b) = 4.25$  GeV,  $M_t^{\text{pole}} = 173.3$  GeV and  $m_\tau^{\text{pole}} = 1.777$  GeV. We have chosen the CKM matrix elements as  $|\mathcal{V}_{tb}| = 0.9993$  and  $|\mathcal{V}_{cb}| = 0.04$ . In order to ensure MFV, for all three generations a common mass parameter  $M_{\tilde{q}_L}$  for the soft SUSY breaking masses of the  $SU(2)$  doublet has to be introduced at the scale  $\mu = \mu_{\text{MFV}}$ , so that the

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<sup>4</sup>Scenarios where 2-body decays at tree level are forbidden for the lightest stop quark and where the loop induced flavour changing decay competes with 3-body decays have been discussed in [22].

up- and down-type squark mass matrices can be simultaneously flavour-diagonal. We work in the framework of a MFV MSSM defined at the GUT scale in terms of a small number of parameters. They are given by common soft SUSY breaking scalar and gaugino mass terms,  $M_0$  and  $M_{1/2}$ , a common SUSY breaking trilinear coupling  $A_0$ , the ratio of the two vacuum expectation values  $\tan\beta$  and the sign of the Higgsino parameter  $\mu$ .

### 5.1 Analysis for $\mu_{\text{MFV}} \approx 10^{16}$ GeV

We first investigate two mSUGRA scenarios with soft-breaking terms at the GUT scale  $M_{\text{GUT}} \approx 10^{16}$  GeV, which is identified with the MFV scale. All soft SUSY breaking parameters are family universal. The boundary conditions at  $\mu_{\text{MFV}} = M_{\text{GUT}}$  are

$$\begin{aligned} (1) \quad & M_0 = 200 \text{ GeV} \quad M_{1/2} = 230 \text{ GeV} \quad A_0 = -920 \text{ GeV} \\ & \tan\beta = 10 \quad \text{sign}(\mu) = + \\ (2) \quad & M_0 = 200 \text{ GeV} \quad M_{1/2} = 230 \text{ GeV} \quad A_0 = -895 \text{ GeV} \\ & \tan\beta = 10 \quad \text{sign}(\mu) = + . \end{aligned} \tag{49}$$

The second scenario has a larger  $\tilde{t}_1 - \tilde{\chi}_1^0$  mass difference compared to scenario (1), whereas the mass difference between  $\tilde{t}_1$  and the lightest chargino  $\tilde{\chi}_1^+$  is smaller. Since the 4-body decays are dominated by the chargino exchange diagram [23], in scenario (2) the 4-body decays should be more important leading to a smaller branching ratio of the flavour changing decay. The GUT scale is given by  $\sim 2.3 \cdot 10^{16}$  GeV. The masses are obtained by RGE evolution from the GUT scale down to the electroweak scale. The running is performed at two-loop order without the inclusion of explicit flavour violation in the squark sector. The obtained masses are

$$\begin{aligned} (1) \quad & m_{\tilde{t}_1} = 104 \text{ GeV} \quad m_{\tilde{\chi}_1^0} = 92 \text{ GeV} \quad m_{\tilde{\chi}_1^+} = 175 \text{ GeV} \\ (2) \quad & m_{\tilde{t}_1} = 130 \text{ GeV} \quad m_{\tilde{\chi}_1^0} = 92 \text{ GeV} \quad m_{\tilde{\chi}_1^+} = 175 \text{ GeV} . \end{aligned} \tag{50}$$

For these scenarios the partial stop decay width into charm and neutralino, calculated with the full one-loop formula, is compared to the approximate result. For the latter, we take  $M_W$  as generic loop particle mass, *cf.* Eq. (47). The widths and form factors are given in Table 1. They have been obtained with the program SUSY-HIT [28], where the full one-loop formula for the flavour changing stop decay has been implemented. As can be inferred from the table, the exact and approximate

$\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$	$\Gamma^{\text{1-loop}}[\text{GeV}]$	$ F_R^{\text{1-loop}} $	$\Gamma^{\text{H/K}}[\text{GeV}]$	$ F_R^{\text{H/K}} $
Scenario(1)	$9.322 \cdot 10^{-10}$	$1.486 \cdot 10^{-4}$	$1.004 \cdot 10^{-9}$	$1.542 \cdot 10^{-4}$
Scenario(2)	$5.862 \cdot 10^{-9}$	$1.460 \cdot 10^{-4}$	$6.446 \cdot 10^{-9}$	$1.531 \cdot 10^{-4}$

Table 1: The partial widths and form factors for the decay  $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$  in two MFV scenarios, calculated with the exact 1-loop formula,  $\Gamma^{\text{1-loop}}$ ,  $F_R^{\text{1-loop}}$ , and with the approximate formula of Ref. [10],  $\Gamma^{\text{H/K}}$ ,  $F_R^{\text{H/K}}$ .

decay width differ by  $\mathcal{O}(10)\%$ . In fact, the finite terms extracted from the one-loop formula turn out to contribute with  $\sim 3 - 5\%$  to  $F_R$ , Eq. (43). This difference leads to the 10% effect in the decay width. The difference in the branching ratio  $BR(\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0)$  calculated in the two approaches

is negligible, however. We note, that in the first scenario the partial width is  $\sim 6$  times smaller than in the second scenario due the smaller  $\tilde{t}_1 - \tilde{\chi}_1^0$  mass difference and hence reduced phase space.

For the calculation of the branching ratios, also the partial width for the  $\tilde{t}_1$  decay into  $u$ -quark and neutralino,  $\tilde{t}_1 \rightarrow u\tilde{\chi}_1^0$ , as well as the 4-body decay width are needed. The former is suppressed by 2 orders of magnitude compared to the  $c\tilde{\chi}_1^0$  final state due to the small CKM matrix element  $|\mathcal{V}_{ub}| \approx 0.003$  which enters quadratically in the decay width. The branching ratios are listed in Table 2. As anticipated, the stop 4-body decay is more important in scenario (2) leading to a change of the branching ratio of interest,  $\text{BR}(\tilde{t}_1 \rightarrow \tilde{\chi}_1^0 c)$ , at the few per-cent level.

branching ratio	$\text{BR}(\tilde{t}_1 \rightarrow \tilde{\chi}_1^0 c)$	$\text{BR}(\tilde{t}_1 \rightarrow \tilde{\chi}_1^0 u)$	$\text{BR}(\tilde{t}_1 \rightarrow \tilde{\chi}_1^0 b f \bar{f}')$
Scenario(1)	0.9944	0.0056	$4.587 \cdot 10^{-5}$
Scenario(2)	0.9443	0.0053	0.0504

Table 2: The  $\tilde{t}_1$  branching ratios for different final states for scenario (1) and (2).

As stated before, the large logarithms in the decay formula should be resummed. To get an estimate of the importance of the resummation effects, the one-loop decay calculated in the framework of MFV is compared to the tree level stop decay into charm and neutralino with flavour off-diagonal elements in the squark mixing matrix. They are the result of the small flavour off-diagonal entries introduced in the soft-breaking terms through RG evolution including the complete flavour structure of the different flavour matrices, from the scale of MFV down to the scale of EWSB. The input parameters for the decay formula are taken from SPheno [25]<sup>5</sup>. In the flavour violating case, denoted by FV in the following, where no flavour-eigenstates exist any more, the lightest up-type squark state  $\tilde{u}_1$  has been identified to correspond to  $\tilde{t}_1$ . The scenarios in Eq. (49) have been chosen such that  $m_{\tilde{u}_1} \approx m_{\tilde{t}_1}$ . The  $\tilde{\chi}_1^0$  and  $\tilde{\chi}_1^\pm$  masses are almost unchanged. The form factor  $F_R^{\text{FV}}$  of the tree level decay is given by the right-handed part of the FCNC  $\tilde{u}_1 - c - \tilde{\chi}_1^0$  coupling<sup>6</sup>,

$$F_R^{\text{FV}} = -i\sqrt{2} \left( \frac{Z_{11}}{6} t_W + \frac{Z_{12}}{2} \right) (\tilde{W}_L)_{\tilde{u}_1 c}, \quad (51)$$

with the squark mixing matrix  $\tilde{W}_L$  defined in Eq. (10). This leads to the partial decay width

$$\Gamma^{\text{FV}}(\tilde{u}_1 \rightarrow c\tilde{\chi}_1^0) = \frac{g^2 m_{\tilde{u}_1}}{16\pi} \left( 1 - \frac{m_{\tilde{\chi}_1^0}^2}{m_{\tilde{u}_1}^2} \right) |F_R^{\text{FV}}|^2. \quad (52)$$

The form factors and partial widths are shown in Table 3. As can be inferred from the table, there is a factor  $\sim 4.4$  between the right-handed form factor calculated at the one-loop level in the MFV framework and the one derived from RG evolution including flavour violation. As expected, resummation effects turn out to be important for a large scale  $\mu_{\text{MFV}} = M_{\text{GUT}}$ <sup>7</sup>. The partial widths,

<sup>5</sup>Thanks to Werner Porod who provided us with the newest SPheno version 3.0.beta56.

<sup>6</sup>The left-handed part is negligibly small for  $m_c = 0$ .

<sup>7</sup>This result is in agreement with the discussion in Ref. [29] where resummation effects in the coupling  $\tilde{t}_1 - c - \tilde{\chi}_1^0$  have been found to be large.

	$ F_R^{1\text{-loop}} $	$ F_R^{\text{FV}} $	$\Gamma^{1\text{-loop}}$ [GeV]	$\Gamma^{\text{FV}}$ [GeV]
Scenario (1)	$1.486 \cdot 10^{-4}$	$3.361 \cdot 10^{-5}$	$9.322 \cdot 10^{-10}$	$4.766 \cdot 10^{-11}$
Scenario (2)	$1.460 \cdot 10^{-4}$	$3.306 \cdot 10^{-5}$	$5.862 \cdot 10^{-9}$	$3.006 \cdot 10^{-10}$

Table 3: The right-handed form factors and partial decay widths of the lightest up-type squark into charm and neutralino for the MFV scenario (1-loop) and the FCNC tree level decay (FV).

which depend quadratically on the right-handed form factor, differ by a factor  $\sim 20$ .

For comparison we have performed the calculation with the decay spectra and mixing angles evaluated by SOFTSUSY. The squark mixing matrix elements agree within  $10^{-2}$  accuracy with the results of SPheno. The mixing matrix element  $(\tilde{W}_L)_{\tilde{u}_1 c}$ , which enters in the form factor  $F_R^{\text{FV}}$  Eq. (51), is  $\mathcal{O}(10^{-4})$  and differs in the two spectrum calculators. The two codes implement the one-loop corrections to the squark mass matrices differently. SOFTSUSY corrects only the flavour-diagonal entries of the squark mass matrices, while SPheno implements a full one-loop calculation, so that differences in the flavour off-diagonal entries are to be expected. For the SOFTSUSY parameter values, this results in a ratio between loop decay and tree level decay of  $\sim 2.7$  for the two scenarios, compared to the ratio  $\sim 4.4$  found with the SPheno parameter values. All in all, the results with both spectrum calculators show the importance of the resummation effects.

Of phenomenological interest are the consequences of these resummation effects on the  $\tilde{t}_1$  branching ratio into the charm plus neutralino final state. To quantify this, the competing 4-body stop decay width, calculated in the FV scenario and including tree level FCNC couplings, is needed. The additional FCNC contributions are expected to be small, however, due to the suppression by CKM matrix elements. As the calculation is not available at present we defer a detailed comparison to a future publication.

## 5.2 Analysis for $\mu_{\text{MFV}} \leq M_{\text{GUT}}$

In the previous section the importance of resummation effects has been discussed. With decreasing  $\mu_{\text{MFV}}$  and hence smaller  $\ln \mu_{\text{MFV}}^2$  on the other hand the non-resummed one-loop MFV result should approach the resummed flavour-violating tree level result. Furthermore, we expect the approximate formula of Ref. [10], which is a good approximation of the exact one-loop MFV result for large scales, to be less good with decreasing MFV scale. In order to verify this behaviour we have chosen scenarios with different  $\mu_{\text{MFV}}$  varied between  $10^3 \text{ GeV} \leq \mu_{\text{MFV}} \leq 10^{16} \text{ GeV}$ . We choose the soft SUSY breaking input parameters<sup>8</sup> in each scenario such that the masses for  $\tilde{t}_1$  and  $\tilde{\chi}_1^0$  remain almost unchanged. Consequently, the differences in the partial decay widths will not be due to phase space effects. Furthermore, the scenarios are constrained by the requirement that  $\tilde{t}_1$  is the NLSP and  $\tilde{\chi}_1^0$  is the LSP, so that the FC decay  $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$  dominates. The relevant particle masses for the different scenarios vary as

$$m_{\tilde{t}_1} = 105 \dots 116 \text{ GeV} \quad \text{and} \quad m_{\tilde{\chi}_1^0} = 92 \dots 104 \text{ GeV} , \quad (53)$$

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<sup>8</sup>In our scenarios  $\tan \beta$  varies between 10 and 20.

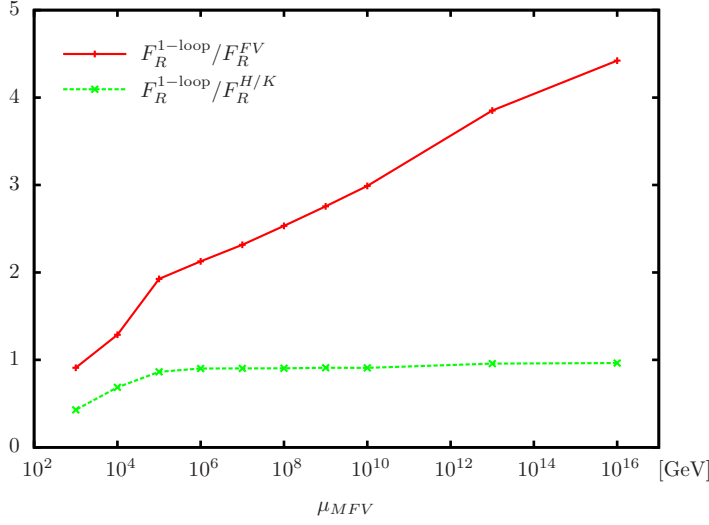


Figure 4: Ratio between the right-handed form factor of the MFV loop decay  $F_R^{1-loop}$  and the form factor of the FV tree level decay  $F_R^{FV}$  (red/full) and ratio between the MFV loop-decay form factor and the approximate form factor  $F_R^{H/K}$  (green/dashed) as function of the MFV scale  $\mu_{MFV}$ .

with the  $\tilde{t}_1 - \tilde{\chi}_1^0$  mass difference ranging between

$$m_{\tilde{t}_1} - m_{\tilde{\chi}_1^0} = 9 \dots 15 \text{ GeV} . \quad (54)$$

We emphasize that the following results are purely illustrative. The various scenarios have not been required to fulfill Dark Matter constraints and/or constraints from electroweak precision data. The main emphasis was to achieve approximately constant masses for the NLSP and LSP.

In Fig. 4 we show, as a function of the MFV scale, the ratio of the non-resummed right-handed form factor  $F_R^{1-loop}$  in the MFV 1-loop decay to  $F_R^{FV}$  in the FV tree level decay as well as the ratio of  $F_R^{1-loop}$  to the approximate form factor  $F_R^{H/K}$ .<sup>9</sup> As can be inferred from the figure, the approximate result reproduces the one-loop result down to low scales. Starting from  $\mu_{MFV} = 10^5$  GeV the finite terms become relevant. At  $\mu_{MFV} = 10^3$  GeV neglecting the finite terms in  $F_R^{H/K}$  leads to a factor  $\sim 2$  between the approximate and the 1-loop form factor. The non-resummed 1-loop result and the resummed tree level result, on the other hand, approach each other with decreasing scale of MFV as expected.

Figure 5 shows the partial widths as functions of  $\mu_{MFV}$  for the approximate MFV decay, for the full MFV 1-loop decay and for the tree level resummed decay. An interesting feature which can be inferred from Fig. 5 is the size of the decay width. It does not only depend on the size of the logarithm but also on the coefficient of the logarithmic term, which is given in terms of the soft SUSY breaking parameters, particle masses and mixing angles, *cf.* Eq. (43). As explained above, for each value of the scale  $\mu_{MFV}$  we have chosen a different set of boundary conditions  $M_0, m_{1/2}, A_0, \tan \beta, \text{sign}(\mu)$  such that the  $\tilde{t}_1$  and  $\tilde{\chi}_1^0$  masses remain approximately unchanged. This leads for each  $\mu_{MFV}$  to a different coefficient of the logarithmic term. For  $\mu_{MFV} = 10^{12}$  GeV *e.g.* the parameter set and resulting masses and mixing angles are such that the coefficient becomes

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<sup>9</sup>Note that the line connecting the different points uniquely serves to guide the eye.



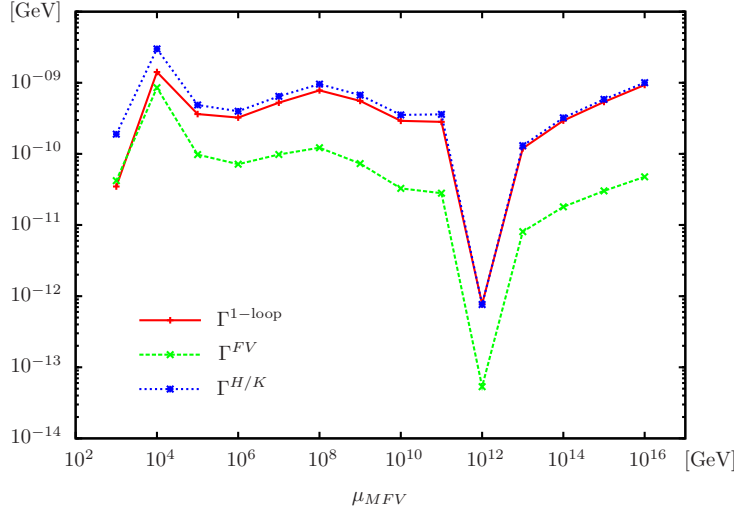


Figure 5: Partial decay width  $\Gamma(\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0)$  calculated assuming MFV,  $\Gamma^{1-loop}$  (red/full), calculated with the approximate formula,  $\Gamma^{H/K}$  (blue/dotted), and calculated at tree level including FV,  $\Gamma^{FV}$  (green/dashed), as function of  $\mu_{MFV}$ .

rather small, so that the partial width is less than  $10^{-12}$  GeV. Due to the large value of  $\mu_{MFV}$  the logarithmic contribution still dominates over the finite terms, however, so that there is good agreement between the 1-loop and approximate result. For small values of  $\mu_{MFV}$  the partial width can be as large as a few  $10^{-11}$  GeV as the factor, which multiplies the logarithm, turns out to be large for the chosen parameter set. The value of the coefficient is also the reason for the kink in Fig. 4 at  $\mu_{MFV} = 10^5$  GeV.

Figure 5 shows, that in accordance with the behaviour of the right-handed form factors, at high scales the 1-loop and the approximate result agree up to the effect of the non-logarithmic terms on the partial width, which is at the 10% level. The 1-loop and the resummed tree level decay agree at low scales where the resummation effects of the large logarithms can be neglected, whereas the deviations are large for high scales. In summary, in order to get correct predictions for the flavour changing light stop decay for large scales of MFV, resummation effects have to be included. To further improve on this decay, the next step is the calculation of the one-loop corrections to the tree level stop decay including the squark mixing matrix elements from RGE evolution. This is deferred to a future publication.

## 6 Summary and Conclusions

In summary, we have calculated the flavour violating decay  $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$  in the framework of Minimal Flavour Violation including also finite terms, which do not depend on the logarithm of the MFV scale  $\mu_{MFV}$ . The one-loop decay has been compared to the approximate result derived earlier by Hikasa and Kobayashi which neglects the subleading terms compared to the large logarithm of  $\mu_{MFV}$ . It has been found that it approximates the complete one-loop result within 10% for large MFV scales. The approximation becomes worse with decreasing scales. The one-loop result, and also the approximate formula, however, do not resum the large logarithms. The resummation is



done by solving the renormalization group equations. Since MFV is not RG-invariant, flavour changing off-diagonal elements are induced in the squark mixing matrices which lead to FCNC couplings at tree level. They can be compared to the effective one-loop coupling in the MFV approach. The resummation effects turn out to be important, so that the one-loop result and the formula by Hikasa and Kobayashi only give an approximate value of the phenomenologically important light stop decay width into charm and neutralino. The next important step to improve the prediction for the light stop decay width will be the calculation of the one-loop corrections to the flavour-violating tree level decay.

## Appendix

### A Couplings

To set up our notation for the couplings, we briefly repeat the chargino and neutralino systems. The chargino mass matrix, in terms of the wino mass parameter  $M_2$ ,  $\mu$  and  $\tan \beta$ , is given by [30]

$$\mathcal{M}_C = \begin{pmatrix} M_2 & \sqrt{2}M_W s_\beta \\ \sqrt{2}M_W c_\beta & \mu \end{pmatrix}, \quad (55)$$

where  $M_W$  denotes the charged  $W$  boson mass and we use  $s_\beta \equiv \sin \beta$ ,  $c_\beta \equiv \cos \beta$ . It is diagonalized by two real matrices  $U$  and  $V$ ,

$$U^* \mathcal{M}_C V^{-1} \rightarrow U = \mathcal{O}_- \quad \text{and} \quad V = \begin{cases} \mathcal{O}_+ & \text{if } \det \mathcal{M}_C > 0 \\ \sigma_3 \mathcal{O}_+ & \text{if } \det \mathcal{M}_C < 0 \end{cases}, \quad (56)$$

with the Pauli matrix  $\sigma_3$  rendering the chargino masses positive.  $\mathcal{O}_\pm$  are rotation matrices with the mixing angles

$$\tan 2\theta_- = \frac{2\sqrt{2}M_W(M_2 c_\beta + \mu s_\beta)}{M_2^2 - \mu^2 - 2M_W^2 c_\beta}, \quad \tan 2\theta_+ = \frac{2\sqrt{2}M_W(M_2 s_\beta + \mu c_\beta)}{M_2^2 - \mu^2 + 2M_W^2 c_\beta}. \quad (57)$$

The two chargino masses read

$$m_{\chi_{1,2}^\pm}^2 = \frac{1}{2} \left\{ M_2^2 + \mu^2 + 2M_W^2 \mp [(M_2^2 - \mu^2)^2 + 4M_W^2(M_W^2 c_{2\beta}^2 + M_2^2 + \mu^2 + 2M_2 \mu s_{2\beta})]^{\frac{1}{2}} \right\}. \quad (58)$$

The four-dimensional neutralino mass matrix in the  $(-i\tilde{B}, -i\tilde{W}_3, \tilde{H}_1^0, \tilde{H}_2^0)$  basis has the form

$$M_N = \begin{pmatrix} M_1 & 0 & -M_Z s_W c_\beta & M_Z s_W s_\beta \\ 0 & M_2 & M_Z c_W c_\beta & -M_Z c_W s_\beta \\ -M_Z s_W c_\beta & M_Z c_W c_\beta & 0 & -\mu \\ M_Z s_W s_\beta & -M_Z c_W s_\beta & -\mu & 0 \end{pmatrix}, \quad (59)$$

with  $c_W^2 = 1 - s_W^2 = M_Z^2/M_Z^2$ . It can be diagonalized analytically [31] with a single matrix  $Z$ .

In the following, we list the couplings in the framework of MFV [30,32–34], which are needed for our calculation. All couplings are normalized to the weak gauge coupling  $g$  if not stated otherwise. Note that all charged couplings involving quarks and/or squarks have to be multiplied with the

CKM matrix element  $\mathcal{V}_{ud}$ , which we have factored out from our definition of the couplings.

- The couplings of charginos and neutralinos to the charged gauge bosons  $W^\pm$ :

$$G_{\tilde{\chi}_i^0 \tilde{\chi}_j^+ W^+}^{L,R} = G_{ij W^+}^{L,R} \quad \text{with} \quad \begin{aligned} G_{ij W^+}^L &= \frac{1}{\sqrt{2}}[-Z_{i4}V_{j2} + \sqrt{2}Z_{i2}V_{j1}] \\ G_{ij W^+}^R &= \frac{1}{\sqrt{2}}[Z_{i3}U_{j2} + \sqrt{2}Z_{i2}U_{j1}] \end{aligned} \quad (60)$$

- The couplings of charginos and neutralinos to charged Higgs and Goldstone bosons<sup>10</sup>,  $H^\pm$ ,  $G^\pm$ :

$$G_{\tilde{\chi}_i^0 \tilde{\chi}_j^+ H^+}^{L,R} = G_{ij H^+}^{L,R} \quad \text{with} \quad \begin{aligned} G_{ij H^+}^L &= c_\beta[Z_{i4}V_{j1} + \frac{1}{\sqrt{2}}(Z_{i2} + \tan \theta_W Z_{i1})V_{j2}] \\ G_{ij H^+}^R &= s_\beta[Z_{i3}U_{j1} - \frac{1}{\sqrt{2}}(Z_{i2} + \tan \theta_W Z_{i1})U_{j2}] \end{aligned} \quad (61)$$

$$G_{\tilde{\chi}_i^0 \tilde{\chi}_j^+ G^+}^{L,R} = G_{ij G^+}^{L,R} \quad \text{with} \quad \begin{aligned} G_{ij G^+}^L &= s_\beta[Z_{i4}V_{j1} + \frac{1}{\sqrt{2}}(Z_{i2} + \tan \theta_W Z_{i1})V_{j2}] \\ G_{ij G^+}^R &= -c_\beta[Z_{i3}U_{j1} - \frac{1}{\sqrt{2}}(Z_{i2} + \tan \theta_W Z_{i1})U_{j2}] \end{aligned} \quad (62)$$

- The couplings between neutralinos, quarks and squarks,  $\tilde{q}_{1,2} - q - \tilde{\chi}_j^0$ :

$$\begin{aligned} \begin{Bmatrix} a_{j1}^q \\ a_{j2}^q \end{Bmatrix} &= -\frac{m_q r_q}{\sqrt{2}M_W} \begin{Bmatrix} s_{\theta_q} \\ c_{\theta_q} \end{Bmatrix} - e_{Lj}^q \begin{Bmatrix} c_{\theta_q} \\ -s_{\theta_q} \end{Bmatrix} \\ \begin{Bmatrix} b_{j1}^q \\ b_{j2}^q \end{Bmatrix} &= -\frac{m_q r_q}{\sqrt{2}M_W} \begin{Bmatrix} c_{\theta_q} \\ -s_{\theta_q} \end{Bmatrix} - e_{Rj}^q \begin{Bmatrix} s_{\theta_q} \\ c_{\theta_q} \end{Bmatrix}, \end{aligned} \quad (63)$$

with  $r_u = Z_{j4}/\sin \beta$  and  $r_d = Z_{j3}/\cos \beta$  for up- and down-type quarks, and

$$\begin{aligned} e_{Lj}^q &= \sqrt{2} [Z_{j1} t_W (Q_q - I_q^3) + Z_{j2} I_q^3] \\ e_{Rj}^q &= -\sqrt{2} Q_q t_W Z_{j1}, \end{aligned} \quad (64)$$

where  $t_W \equiv \tan \theta_W$ .

- The couplings between charginos, quarks and squarks,  $\tilde{q}_{1,2} - q' - \tilde{\chi}_j^+$ , for up- and down-type (s)quarks read

$$\begin{aligned} \begin{Bmatrix} a_{j1}^{\tilde{u}d} \\ a_{j2}^{\tilde{u}d} \end{Bmatrix} &= V_{j1} \begin{Bmatrix} -c_{\theta_u} \\ s_{\theta_u} \end{Bmatrix} + \frac{m_u V_{j2}}{\sqrt{2}M_W s_\beta} \begin{Bmatrix} s_{\theta_u} \\ c_{\theta_u} \end{Bmatrix} \\ \begin{Bmatrix} b_{j1}^{\tilde{u}d} \\ b_{j2}^{\tilde{u}d} \end{Bmatrix} &= \frac{m_d U_{j2}}{\sqrt{2}M_W c_\beta} \begin{Bmatrix} c_{\theta_u} \\ -s_{\theta_u} \end{Bmatrix} \end{aligned} \quad (65)$$

$$\begin{aligned} \begin{Bmatrix} a_{j1}^{\tilde{d}u} \\ a_{j2}^{\tilde{d}u} \end{Bmatrix} &= U_{j1} \begin{Bmatrix} -c_{\theta_d} \\ s_{\theta_d} \end{Bmatrix} + \frac{m_d U_{j2}}{\sqrt{2}M_W c_\beta} \begin{Bmatrix} s_{\theta_d} \\ c_{\theta_d} \end{Bmatrix} \\ \begin{Bmatrix} b_{j1}^{\tilde{d}u} \\ b_{j2}^{\tilde{d}u} \end{Bmatrix} &= \frac{m_u V_{j2}}{\sqrt{2}M_W s_\beta} \begin{Bmatrix} c_{\theta_d} \\ -s_{\theta_d} \end{Bmatrix} \end{aligned} \quad (66)$$

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<sup>10</sup>We work in the Feynman gauge.

- The couplings of the  $W^\pm$  gauge bosons, the charged Higgs and Goldstone bosons to quarks are given by

$$v_W^q = -a_W^q = -\frac{1}{2\sqrt{2}} \quad (67)$$

$$v_{H^+}^q = \frac{m_d \tan \beta + m_u \cot \beta}{2\sqrt{2}M_W}, \quad a_{H^+}^q = \frac{m_d \tan \beta - m_u \cot \beta}{2\sqrt{2}M_W} \quad (68)$$

$$v_{G^+}^q = \frac{-m_d + m_u}{2\sqrt{2}M_W}, \quad a_{G^+}^q = -\frac{m_d + m_u}{2\sqrt{2}M_W}. \quad (69)$$

- The couplings of the  $W^\pm$  gauge bosons and charged Higgs and Goldstone bosons to squarks are given by

$$G_{W^+\tilde{u}_i\tilde{d}_j} = -\frac{1}{\sqrt{2}}C_{W^+ij} \quad \text{with}$$

$$\begin{aligned} C_{W^+11} &= c_{\theta_u}c_{\theta_d} & C_{W^+12} &= -c_{\theta_u}s_{\theta_d} \\ C_{W^+21} &= -s_{\theta_u}c_{\theta_d} & C_{W^+22} &= s_{\theta_u}s_{\theta_d} \end{aligned} \quad (70)$$

$$G_{H^+\tilde{u}_i\tilde{d}_j} = C_{H^+ij} \quad \text{with}$$

$$\begin{aligned} C_{H^+11} &= c_{\theta_u}c_{\theta_d}s_{11} + s_{\theta_u}s_{\theta_d}s_{22} + c_{\theta_u}s_{\theta_d}s_{12} + s_{\theta_u}c_{\theta_d}s_{21} \\ C_{H^+22} &= s_{\theta_u}s_{\theta_d}s_{11} + c_{\theta_u}c_{\theta_d}s_{22} - s_{\theta_u}c_{\theta_d}s_{12} - c_{\theta_u}s_{\theta_d}s_{21} \\ C_{H^+12} &= -c_{\theta_u}s_{\theta_d}s_{11} + s_{\theta_u}c_{\theta_d}s_{22} + c_{\theta_u}c_{\theta_d}s_{12} - s_{\theta_u}s_{\theta_d}s_{21} \\ C_{H^+21} &= -s_{\theta_u}c_{\theta_d}s_{11} + c_{\theta_u}s_{\theta_d}s_{22} - s_{\theta_u}s_{\theta_d}s_{12} + c_{\theta_u}c_{\theta_d}s_{21}, \end{aligned} \quad (71)$$

where

$$\begin{aligned} s_{11} &= -\frac{M_W}{\sqrt{2}} \left( \sin 2\beta - \frac{m_d^2 \tan \beta + m_u^2 \cot \beta}{M_W^2} \right) \\ s_{22} &= \frac{m_u m_d}{\sqrt{2}M_W} (\tan \beta + \cot \beta) \\ s_{12} &= \frac{m_d}{\sqrt{2}M_W} (\mu + A_d \tan \beta) \\ s_{21} &= \frac{m_u}{\sqrt{2}M_W} (\mu + A_u \cot \beta). \end{aligned} \quad (72)$$

And for the Goldstone couplings we have

$$G_{G^+\tilde{u}_i\tilde{d}_j} = C_{G^+ij} \quad \text{with}$$

$$\begin{aligned} C_{G^+11} &= c_{\theta_u}c_{\theta_d}a_{LL} + c_{\theta_u}s_{\theta_d}a_{LR} + s_{\theta_u}c_{\theta_d}a_{RL} \\ C_{G^+22} &= s_{\theta_u}s_{\theta_d}a_{LL} - s_{\theta_u}c_{\theta_d}a_{LR} - c_{\theta_u}s_{\theta_d}a_{RL} \\ C_{G^+12} &= -c_{\theta_u}s_{\theta_d}a_{LL} + c_{\theta_u}c_{\theta_d}a_{LR} - s_{\theta_u}s_{\theta_d}a_{RL} \\ C_{G^+21} &= -s_{\theta_u}c_{\theta_d}a_{LL} - s_{\theta_u}s_{\theta_d}a_{LR} + c_{\theta_u}c_{\theta_d}a_{RL}, \end{aligned} \quad (73)$$

and

$$\begin{aligned}
a_{LL} &= \frac{M_W}{\sqrt{2}} \left( \cos 2\beta + \frac{m_u^2 - m_d^2}{M_W^2} \right) \\
a_{LR} &= \frac{m_d}{\sqrt{2}M_W} (\mu \tan \beta - A_d) \\
a_{RL} &= -\frac{m_u}{\sqrt{2}M_W} (\mu \cot \beta - A_u) .
\end{aligned} \tag{74}$$

• For our calculation we also need the 4-squark coupling between stop, scharm and two *identical* down-type squarks. With the generation index  $k = 1, 2, 3$  and the index  $l = 1, 2$  denoting the two down-type squark eigenstates, it reads in terms of  $g^2$

$$G_{\tilde{t}_i \tilde{c}_j \tilde{d}_{kl} \tilde{d}_{kl}} = C_{ij\tilde{d}_{kl}\tilde{d}_{kl}} \quad \text{with} \quad C_{ij\tilde{d}_{kl}\tilde{d}_{kl}} = -\frac{m_t m_c}{2M_W^2 s_\beta^2} \mathcal{P}_{ijll} - \frac{m_{d_k}^2}{2M_W^2 c_\beta^2} \mathcal{P}'_{ijll} - \frac{1}{2} \mathcal{P}''_{ijll} \tag{75}$$

and

$ijll$	$\mathcal{P}_{ijll}$	$\mathcal{P}'_{ijll}$	$\mathcal{P}''_{ijll}$
1111	$s_{\theta_t} s_{\theta_c} c_{\theta_{d_k}}^2$	$c_{\theta_t} c_{\theta_c} s_{\theta_{d_k}}^2$	$c_{\theta_c} c_{\theta_t} c_{\theta_{d_k}}^2$
2211	$c_{\theta_t} c_{\theta_c} c_{\theta_{d_k}}^2$	$s_{\theta_c} s_{\theta_t} s_{\theta_{d_k}}^2$	$s_{\theta_c} s_{\theta_t} c_{\theta_{d_k}}^2$
1211	$s_{\theta_t} c_{\theta_c} c_{\theta_{d_k}}^2$	$-c_{\theta_t} s_{\theta_c} s_{\theta_{d_k}}^2$	$-c_{\theta_t} s_{\theta_c} c_{\theta_{d_k}}^2$
2111	$c_{\theta_t} s_{\theta_c} c_{\theta_{d_k}}^2$	$-c_{\theta_c} s_{\theta_t} s_{\theta_{d_k}}^2$	$-c_{\theta_c} s_{\theta_t} c_{\theta_{d_k}}^2$

(76)

The couplings with two  $\tilde{d}_{k2}$  squarks are obtained from those with the  $\tilde{d}_{k1}$  squarks by interchanging  $c_{\theta_{d_k}} \leftrightarrow s_{\theta_{d_k}}$ .

• Finally, the couplings between stop, scharm and two charged Higgs bosons or two charged Goldstone bosons in terms of  $g^2$  are given by

$$G_{\tilde{t}_i \tilde{c}_j H^+ H^+} = -\frac{m_b^2 \tan \beta^2}{2M_W^2} Q_{ijH^+ H^+} \tag{77}$$

and

$$G_{\tilde{t}_i \tilde{c}_j G^+ G^+} = -\frac{m_b^2}{2M_W^2} Q_{ijH^+ H^+} \tag{78}$$

where

$$\begin{aligned}
Q_{11H^+ H^+} &= c_{\theta_t} c_{\theta_c} & Q_{22H^+ H^+} &= s_{\theta_t} s_{\theta_c} \\
Q_{12H^+ H^+} &= -c_{\theta_t} s_{\theta_c} & Q_{21H^+ H^+} &= -s_{\theta_t} c_{\theta_c}
\end{aligned} . \tag{79}$$

## B1 Squark self-energy contributions

In this Appendix we give the result for the self-energy  $\tilde{\Sigma}_{\tilde{c}_L \tilde{t}_1}$  in terms of the couplings defined in Appendix A. The self-energy receives contributions from the various diagrams in Fig.2, *i.e.*

$$\tilde{\Sigma}_{\tilde{c}_L \tilde{t}_1} = \Sigma_{\tilde{c}_L \tilde{t}_1}^{\tilde{\chi}^+ d} + \Sigma_{\tilde{c}_L \tilde{t}_1}^{H^+ \tilde{d}} + \Sigma_{\tilde{c}_L \tilde{t}_1}^{G^+ \tilde{d}} + \Sigma_{\tilde{c}_L \tilde{t}_1}^{W^+ \tilde{d}} + \Sigma_{\tilde{c}_L \tilde{t}_1}^{H^+ G^+} + \Sigma_{\tilde{c}_L \tilde{t}_1}^{4\tilde{q}} . \tag{80}$$

Note that the self-energy for  $\tilde{t}_1$  and  $\tilde{c}_R$  external legs is zero for vanishing  $c$  quark mass,

$$\tilde{\Sigma}_{\tilde{c}_R \tilde{t}_1} |_{m_c=0} = 0. \quad (81)$$

We have for

$$\begin{aligned} \Sigma_{\tilde{c}_L \tilde{t}_1}^{\tilde{X}^+ d}(m_{\tilde{t}_1}^2) &= g^2 \sum_{j=1,2} \sum_{k=1,\dots,3} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* (-2) \left\{ \frac{1}{2} (a_{j1}^{\tilde{c}d_k} a_{j1}^{\tilde{t}d_k} + b_{j1}^{\tilde{c}d_k} b_{j1}^{\tilde{t}d_k}) [A_0(m_{\tilde{X}_j^+}^2) + A_0(m_{d_k}^2)] \right. \\ &+ (m_{\tilde{X}_j^+}^2 - m_{\tilde{t}_1}^2 + m_{d_k}^2) B_0(m_{\tilde{t}_1}^2, m_{\tilde{X}_j^+}, m_{d_k})] \\ &+ (a_{j1}^{\tilde{c}d_k} b_{j1}^{\tilde{t}d_k} + a_{j1}^{\tilde{t}d_k} b_{j1}^{\tilde{c}d_k}) m_{\tilde{X}_j^+} m_{d_k} B_0(m_{\tilde{t}_1}^2, m_{\tilde{X}_j^+}, m_{d_k}) \left. \right\}, \end{aligned} \quad (82)$$

with  $a_{j1}^{\tilde{c}d_k}$  etc. given in Eq. (65). The sum is taken over the chargino eigenstates and the three quark generations. The scalar one-loop one- and two-point integrals  $A_0(m)$  and  $B_0(p^2, m_1, m_2)$  are defined as [35]

$$A_0(m) = -i \bar{\mu}^{4-n} \int \frac{d^n k}{(2\pi)^n} \frac{1}{k^2 - m^2} \quad (83)$$

$$B_0(p^2, m_1, m_2) = -i \bar{\mu}^{4-n} \int \frac{d^n k}{(2\pi)^n} \frac{1}{(k^2 - m_1^2)[(k+p)^2 - m_2^2]}. \quad (84)$$

Note the suppression by the CKM matrix elements  $\mathcal{V}_{cd_k}, \mathcal{V}_{td_k}^*$ . We find for  $\Sigma_{\tilde{c}_L \tilde{t}_1}^{H^+ \tilde{d}}$ :

$$\Sigma_{\tilde{c}_L \tilde{t}_1}^{H^+ \tilde{d}}(m_{\tilde{t}_1}^2) = g^2 \sum_{k=1,\dots,3} \sum_{i=1,2} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* G_{H^+ \tilde{t}_1 \tilde{d}_{ki}} G_{H^+ \tilde{c}_L \tilde{d}_{ki}} B_0(m_{\tilde{t}_1}^2, m_{H^+}, m_{\tilde{d}_{ki}}). \quad (85)$$

The sum is to be taken over the three squark generations  $k$  and the two squark mass eigenstates  $i$ . The Goldstone contribution reads

$$\Sigma_{\tilde{c}_L \tilde{t}_1}^{G^+ \tilde{d}}(m_{\tilde{t}_1}^2) = g^2 \sum_{k=1,\dots,3} \sum_{i=1,2} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* G_{G^+ \tilde{t}_1 \tilde{d}_{ki}} G_{G^+ \tilde{c}_L \tilde{d}_{ki}} B_0(m_{\tilde{t}_1}^2, M_W, m_{\tilde{d}_{ki}}), \quad (86)$$

and the self-energy involving the  $W$  boson

$$\begin{aligned} \Sigma_{\tilde{c}_L \tilde{t}_1}^{W^+ \tilde{d}}(m_{\tilde{t}_1}^2) &= g^2 \sum_{k=1,\dots,3} \sum_{i=1,2} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* G_{W^+ \tilde{t}_1 \tilde{d}_{ki}} G_{W^+ \tilde{c}_L \tilde{d}_{ki}} \left\{ -2A_0(M_W^2) + A_0(m_{\tilde{d}_{ki}}^2) \right. \\ &- (2m_{\tilde{t}_1}^2 + 2m_{\tilde{d}_{ki}}^2 - M_W^2) B_0(m_{\tilde{t}_1}^2, M_W, m_{\tilde{d}_{ki}}) \left. \right\}. \end{aligned} \quad (87)$$

Finally, we have the tadpole contributions from the charged Higgs and Goldstone boson loop,

$$\Sigma_{\tilde{c}_L \tilde{t}_1}^{H^+ G^+}(m_{\tilde{t}_1}^2) = g^2 (-1) \mathcal{V}_{cb} \mathcal{V}_{tb}^* [G_{\tilde{t}_1 \tilde{c}_L H^+ H^+} A_0(m_{H^+}^2) + G_{\tilde{t}_1 \tilde{c}_L G^+ G^+} A_0(M_W^2)], \quad (88)$$

and the one from the 4-squark vertex

$$\Sigma_{\tilde{c}_L \tilde{t}_1}^{4\tilde{q}}(m_{\tilde{t}_1}^2) = g^2 \sum_{k=1,\dots,3} \sum_{i=1,2} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* G_{\tilde{t}_1 \tilde{c}_L \tilde{d}_{ki} \tilde{d}_{ki}} A_0(m_{\tilde{d}_{ki}}^2). \quad (89)$$

## B2 Quark self-energy

According to the structure of the quark self-energy given in Eq. (22) we find for the left-chiral contribution

$$\begin{aligned}\Sigma_{tc}^L(p^2) &= g^2 \left\{ \mathcal{V}_{cb} \mathcal{V}_{tb}^* \frac{m_b^2}{2M_W^2} [-t_\beta^2 B_1(p^2, m_b, m_{H^+}) - B_1(p^2, m_b, M_W)] \right. \\ &\quad + \sum_{j=1,2} \sum_{k=1,\dots,3} \sum_{i=1,2} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* (-a_{ji}^{\tilde{d}_k c} a_{ji}^{\tilde{d}_k t}) B_1(p^2, m_{\tilde{\chi}_j^+}, m_{\tilde{d}_{ki}}) \\ &\quad \left. + \sum_{k=1,\dots,3} \frac{1}{2} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* [-1 - 2B_1(p^2, m_{d_k}, M_W)] \right\},\end{aligned}\quad (90)$$

with  $a_{ji}^{\tilde{d}_k c}, a_{ji}^{\tilde{d}_k t}$  given in Eq. (66). The sums are taken over all possible chargino states ( $j = 1, 2$ ), the three quark and squark generations ( $k = 1, 2, 3$ ) and the two squark mass eigenstates ( $i = 1, 2$ ). Furthermore,  $B_1$  in terms of the scalar one- and two-point functions is given by

$$B_1(p^2, m_0, m_1) = \frac{1}{2p^2} [A_0(m_0) - A_0(m_1) - (p^2 - m_1^2 + m_0^2) B_0(p^2, m_0, m_1)] . \quad (91)$$

The right-chiral contribution reads

$$\begin{aligned}\Sigma_{tc}^R(p^2) &= g^2 \left\{ \sum_{k=1,\dots,3} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* \frac{m_c m_t}{2M_W^2} [-\cot_\beta^2 B_1(p^2, m_{d_k}, m_{H^+}) - B_1(p^2, m_{d_k}, M_W)] \right. \\ &\quad \left. + \sum_{j=1,2} \sum_{k=1,2,3} \sum_{i=1,2} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* (-b_{ji}^{\tilde{d}_k c} b_{ji}^{\tilde{d}_k t}) B_1(p^2, m_{\tilde{\chi}_j^+}, m_{\tilde{d}_{ki}}) \right\},\end{aligned}\quad (92)$$

with  $b_{ji}^{\tilde{d}_k c}, b_{ji}^{\tilde{d}_k t}$  defined in Eq. (66). It vanishes for  $m_c = 0$ . The left-chiral scalar contribution can be cast into the form

$$\begin{aligned}\Sigma_{tc}^{Ls}(p^2) &= g^2 \left\{ \mathcal{V}_{cb} \mathcal{V}_{tb}^* \frac{m_b^2 m_c}{2M_W^2} [B_0(p^2, m_{H^+}, m_b) - B_0(p^2, M_W, m_b)] \right. \\ &\quad \left. + \sum_{j=1,2} \sum_{k=1,\dots,3} \sum_{i=1,2} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* m_{\tilde{\chi}_j^+} a_{ji}^{\tilde{d}_k t} b_{ji}^{\tilde{d}_k c} B_0(p^2, m_{\tilde{d}_{ki}}, m_{\tilde{\chi}_j^+}) \right\},\end{aligned}\quad (93)$$

which also vanishes for zero charm quark mass. For the right-chiral scalar contribution we find

$$\begin{aligned}\Sigma_{tc}^{Rs}(p^2) &= g^2 \left\{ \mathcal{V}_{cb} \mathcal{V}_{tb}^* \frac{m_b^2 m_t}{2M_W^2} [B_0(p^2, m_{H^+}, m_b) - B_0(p^2, M_W, m_b)] \right. \\ &\quad \left. + \sum_{j=1,2} \sum_{k=1,\dots,3} \sum_{i=1,2} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* m_{\tilde{\chi}_j^+} b_{ji}^{\tilde{d}_k t} a_{ji}^{\tilde{d}_k c} B_0(p^2, m_{\tilde{d}_{ki}}, m_{\tilde{\chi}_j^+}) \right\}.\end{aligned}\quad (94)$$

Note that in case of real two-point functions we have  $\Sigma_{tc}^{Rs} = \Sigma_{ct}^{Ls*}$ .

## B3 Vertex correction

The vertex contributions to the left-chiral form factor  $F_L^v$  vanish for  $m_c = 0$ . For the right-chiral form factor  $F_R^v$  they are given by the various right-chiral contributions from the vertex correction graphs depicted in Fig.2

$$\begin{aligned}gF_R^v(m_{\tilde{t}_1}^2) &= i [\Gamma_{\tilde{\chi}^+ \tilde{d} d} + \Gamma_{\tilde{\chi}^+ H^+ d} + \Gamma_{\tilde{\chi}^+ G^+ d} + \Gamma_{\tilde{\chi}^+ W^+ d} + \Gamma_{\tilde{d} H^+ d} + \Gamma_{\tilde{d} G^+ d} + \Gamma_{\tilde{\chi}^+ \tilde{d} H^+} + \Gamma_{\tilde{\chi}^+ \tilde{d} G^+} \\ &\quad + \Gamma_{\tilde{d} W^+} + \Gamma_{\tilde{d} \tilde{\chi}^+ W^+}] (m_{\tilde{t}_1}^2) .\end{aligned}\quad (95)$$

We have for  $\Gamma_{\tilde{\chi}^+ \tilde{d} d}$ ,

$$\begin{aligned}
\Gamma_{\tilde{\chi}^+ \tilde{d} d}(m_{\tilde{t}_1}^2) &= -g^3 \sum_{j=1,2} \sum_{k=1\dots 3} \sum_{i=1,2} \mathcal{V}_{cdk} \mathcal{V}_{tdk}^* \left\{ [c_{1ijk} m_{\tilde{\chi}_1^0} + c_{2ijk} m_{\tilde{t}_1}^2] B_0(m_{\tilde{t}_1}^2, m_{d_k}, m_{\tilde{\chi}_j^+}) \right. \\
&\quad - [c_{1ijk} m_{\tilde{\chi}_1^0} + c_{2ijk} m_{\tilde{\chi}_1^0}^2] B_0(m_{\tilde{\chi}_1^0}^2, m_{\tilde{d}_{ki}}, m_{d_k}) + [c_{1ijk} m_{\tilde{\chi}_1^0} (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2 + m_{\tilde{d}_{ki}}^2 - m_{\tilde{\chi}_j^+}^2) \\
&\quad - c_{2ijk} (m_{\tilde{\chi}_j^+}^2 m_{\tilde{\chi}_1^0}^2 - m_{\tilde{t}_1}^2 m_{\tilde{d}_{ki}}^2) + (-c_{3ijk} m_{\tilde{\chi}_1^0} m_{d_k} + c_{4ijk} m_{\tilde{\chi}_j^+} m_{d_k}) (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2)] \\
&\quad \left. C_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_1^0}^2, 0, m_{\tilde{\chi}_j^+}, m_{d_k}, m_{\tilde{d}_{ki}}) \right\} / (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2), \tag{96}
\end{aligned}$$

where the sum over all possible chargino eigenstates  $\tilde{\chi}_j^+$  ( $j = 1, 2$ ), all three generations of down type quarks and squarks ( $k = 1, 2, 3$ ) as well as the two squark mass eigenstates ( $i = 1, 2$ ) has to be taken. We have introduced the abbreviations

$$\begin{aligned}
c_{1ijk} &= a_{ji}^{\tilde{d}_k c} a_{j1}^{\tilde{t}_k} a_{1i}^{d_k} m_{\tilde{\chi}_j^+} + a_{ji}^{\tilde{d}_k c} b_{j1}^{\tilde{t}_k} a_{1i}^{d_k} m_{d_k} & c_{2ijk} &= a_{ji}^{\tilde{d}_k c} b_{j1}^{\tilde{t}_k} b_{1i}^{d_k} \\
c_{3ijk} &= a_{ji}^{\tilde{d}_k c} b_{j1}^{\tilde{t}_k} a_{1i}^{d_k} & c_{4ijk} &= a_{ji}^{\tilde{d}_k c} a_{j1}^{\tilde{t}_k} b_{1i}^{d_k}, \tag{97}
\end{aligned}$$

with the various couplings defined in Eqs. (63,65,66). The scalar one-loop 3-point function is given by

$$\begin{aligned}
C_0(p_1^2, p_2^2, (p_1 + p_2)^2, m_1, m_2, m_3) &= \\
&= -i \bar{\mu}^{4-n} \int \frac{d^n k}{(2\pi)^n} \frac{1}{(k^2 - m_1^2)[(k + p_1)^2 - m_2^2][(k + p_1 + p_2)^2 - m_3^2]}. \tag{98}
\end{aligned}$$

We find

$$\begin{aligned}
\Gamma_{\tilde{\chi}^+ H^+ d}(m_{\tilde{t}_1}^2) &= g^3 \sum_{j=1,2} \mathcal{V}_{cb} \mathcal{V}_{tb}^* \frac{m_b \tan \beta}{\sqrt{2} M_W} \left\{ [c_{1j} m_{\tilde{\chi}_1^0} + c_{2j} m_{\tilde{t}_1}^2] B_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_j^+}, m_b) \right. \\
&\quad - [c_{1j} m_{\tilde{\chi}_1^0} + c_{2j} m_{\tilde{\chi}_1^0}^2] B_0(m_{\tilde{\chi}_1^0}^2, m_{H^+}, m_{\tilde{\chi}_j^+}) + [c_{1j} m_{\tilde{\chi}_1^0} (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2 - m_b^2 + m_{H^+}^2) \\
&\quad + c_{2j} (m_{\tilde{t}_1}^2 m_{H^+}^2 - m_{\tilde{\chi}_1^0}^2 m_b^2) + m_{\tilde{\chi}_j^+} (c_{3j} m_b - c_{4j} m_{\tilde{\chi}_1^0}) (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2)] \\
&\quad \left. C_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_1^0}^2, 0, m_b, m_{\tilde{\chi}_j^+}, m_{H^+}) \right\} / (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2), \tag{99}
\end{aligned}$$

with

$$\begin{aligned}
c_{1j} &= a_{j1}^{\tilde{t} b} G_{1jH^+}^L m_b + b_{j1}^{\tilde{t} b} G_{1j4}^L m_{\tilde{\chi}_j^+} & c_{2j} &= b_{j1}^{\tilde{t} b} G_{1jH^+}^R \\
c_{3j} &= a_{j1}^{\tilde{t} b} G_{1jH^+}^R & c_{4j} &= b_{j1}^{\tilde{t} b} G_{1jH^+}^L.
\end{aligned}$$

Note that we set the down and strange quark mass to zero,  $m_d = m_s = 0$ . We have

$$\begin{aligned}
\Gamma_{\tilde{\chi}^+ G^+ d}(m_{\tilde{t}_1}^2) &= g^3 \sum_{j=1,2} \mathcal{V}_{cb} \mathcal{V}_{tb}^* \frac{-m_b}{\sqrt{2} M_W} \left\{ [c_{5j} m_{\tilde{\chi}_1^0} + c_{6j} m_{\tilde{t}_1}^2] B_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_j^+}, m_b) \right. \\
&\quad - [c_{5j} m_{\tilde{\chi}_1^0} + c_{6j} m_{\tilde{\chi}_1^0}^2] B_0(m_{\tilde{\chi}_1^0}^2, M_W, m_{\tilde{\chi}_j^+}) + [c_{5j} m_{\tilde{\chi}_1^0} (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2 - m_b^2 + M_W^2) \\
&\quad + c_{6j} (m_{\tilde{t}_1}^2 M_W^2 - m_{\tilde{\chi}_1^0}^2 m_b^2) + m_{\tilde{\chi}_j^+} (c_{7j} m_b - c_{8j} m_{\tilde{\chi}_1^0}) (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2)] \\
&\quad \left. C_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_1^0}^2, 0, m_b, m_{\tilde{\chi}_j^+}, M_W) \right\} / (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2), \tag{100}
\end{aligned}$$

where

$$\begin{aligned} c_{5j} &= a_{j1}^{\tilde{t}b} G_{1jG^+}^L m_b + b_{j1}^{\tilde{t}b} G_{1jG^+}^L m_{\tilde{\chi}_j^+} & c_{6j} &= b_{j1}^{\tilde{t}b} G_{1jG^+}^R \\ c_{7j} &= a_{j1}^{\tilde{t}b} G_{1jG^+}^R & c_{8j} &= b_{j1}^{\tilde{t}b} G_{1jG^+}^L . \end{aligned}$$

And

$$\begin{aligned} \Gamma_{\tilde{\chi}^+W+d}(m_{\tilde{t}_1}^2) &= g^3 \sum_{j=1,2} \sum_{k=1\dots 3} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* \frac{1}{\sqrt{2}} \left\{ -2c_{1jk} m_{\tilde{\chi}_1^0} B_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_j^+}, m_{d_k}) \right. \\ &+ 2[c_{1jk} m_{\tilde{\chi}_1^0} + c_{2jk}(m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2)] B_0(m_{\tilde{\chi}_1^0}^2, M_W, m_{\tilde{\chi}_j^+}) + 2c_{2jk}(m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) \\ &B_0(0, M_W, m_{d_k}) + [2c_{2jk}(m_{d_k}^2 + m_{\tilde{\chi}_j^+}^2 - m_{\tilde{t}_1}^2)(m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) + 2m_{\tilde{\chi}_1^0} c_{1jk} \\ &(m_{\tilde{\chi}_1^0}^2 - m_{\tilde{t}_1}^2 + m_{d_k}^2 - M_W^2) + 2m_{\tilde{\chi}_j^+}(c_{3jk} m_{\tilde{\chi}_1^0} + 2c_{4jk} m_{d_k})(m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2)] \\ &\left. C_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_1^0}^2, 0, m_{d_k}, m_{\tilde{\chi}_j^+}, M_W) - 2c_{2jk}(m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) \right\} / (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) , \quad (101) \end{aligned}$$

with

$$\begin{aligned} c_{1jk} &= b_{j1}^{\tilde{t}d_k} G_{1jW^+m_{d_k}}^R + a_{j1}^{\tilde{t}d_k} G_{1jW^+m_{\tilde{\chi}_j^+}}^R & c_{2jk} &= a_{j1}^{\tilde{t}d_k} G_{1jW^+}^L \\ c_{3jk} &= a_{j1}^{\tilde{t}d_k} G_{1jW^+}^R & c_{4jk} &= b_{j1}^{\tilde{t}d_k} G_{1jW^+}^L . \end{aligned} \quad (102)$$

Next

$$\begin{aligned} \Gamma_{\tilde{d}H+d}(m_{\tilde{t}_1}^2) &= g^3 \sum_{i=1,2} \mathcal{V}_{cb} \mathcal{V}_{tb}^* \frac{m_b \tan \beta G_{H^+\tilde{t}_1\tilde{b}_i}}{\sqrt{2} M_W} \left\{ b_{1i}^b m_{\tilde{\chi}_1^0} [B_0(m_{\tilde{t}_1}^2, m_{\tilde{b}_i}, m_{H^+}) - B_0(m_{\tilde{\chi}_1^0}^2, m_b, m_{\tilde{b}_i})] \right. \\ &- [a_{1i}^b m_b(m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) - b_{1i}^b m_{\tilde{\chi}_1^0}(m_b^2 - m_{H^+}^2)] C_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_1^0}^2, 0, m_{H^+}, m_{\tilde{b}_i}, m_b) \left. \right\} \\ &/ (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) \quad (103) \end{aligned}$$

and

$$\begin{aligned} \Gamma_{\tilde{d}G+d}(m_{\tilde{t}_1}^2) &= g^3 \sum_{i=1,2} \mathcal{V}_{cb} \mathcal{V}_{tb}^* \frac{-m_b G_{G^+\tilde{t}_1\tilde{b}_i}}{\sqrt{2} M_W} \left\{ b_{1i}^b m_{\tilde{\chi}_1^0} [B_0(m_{\tilde{t}_1}^2, m_{\tilde{b}_i}, M_W) - B_0(m_{\tilde{\chi}_1^0}^2, m_b, m_{\tilde{b}_i})] \right. \\ &- [a_{1i}^b m_b(m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) - b_{1i}^b m_{\tilde{\chi}_1^0}(m_b^2 - M_W^2)] C_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_1^0}^2, 0, M_W, m_{\tilde{b}_i}, m_b) \left. \right\} \\ &/ (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) . \quad (104) \end{aligned}$$

Furthermore,

$$\begin{aligned} \Gamma_{\tilde{\chi}^+\tilde{d}H^+}(m_{\tilde{t}_1}^2) &= g^3 \sum_{j=1,2} \sum_{k=1\dots 3} \sum_{i=1,2} G_{H^+\tilde{t}_1\tilde{d}_{ki}} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* \left\{ c_{5ijk} m_{\tilde{\chi}_1^0} [-B_0(m_{\tilde{t}_1}^2, m_{H^+}, m_{\tilde{d}_{ki}}) \right. \\ &+ B_0(m_{\tilde{\chi}_1^0}^2, m_{\tilde{\chi}_j^+}, m_{H^+})] + [c_{6ijk} m_{\tilde{\chi}_j^+}(m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) + c_{5ijk} m_{\tilde{\chi}_1^0}(m_{\tilde{d}_{ki}}^2 - m_{\tilde{\chi}_j^+}^2)] \\ &\left. C_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_1^0}^2, 0, m_{\tilde{d}_{ki}}, m_{H^+}, m_{\tilde{\chi}_j^+}) \right\} / (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) \quad (105) \end{aligned}$$

with

$$c_{5ijk} = G_{1jH^+}^R a_{ji}^{\tilde{d}_{ki}c} \quad c_{6ijk} = G_{1jH^+}^L a_{ji}^{\tilde{d}_{ki}c} , \quad (106)$$



and

$$\begin{aligned}
\Gamma_{\tilde{\chi}^+ \tilde{d} G^+}(m_{\tilde{t}_1}^2) &= g^3 \sum_{j=1,2} \sum_{k=1\dots 3} \sum_{i=1,2} G_{G^+ \tilde{t}_1 \tilde{d}_{ki}} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* \left\{ c_{7ijk} m_{\tilde{\chi}_1^0} [-B_0(m_{\tilde{t}_1}^2, M_W, m_{\tilde{d}_{ki}}) \right. \\
&+ B_0(m_{\tilde{\chi}_1^0}^2, m_{\tilde{\chi}_j^+}, M_W)] + [c_{8ijk} m_{\tilde{\chi}_j^+} (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) + c_{7ijk} m_{\tilde{\chi}_1^0} (m_{\tilde{d}_{ki}}^2 - m_{\tilde{\chi}_j^+}^2)] \\
&\left. C_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_1^0}^2, 0, m_{\tilde{d}_{ki}}, M_W, m_{\tilde{\chi}_j^+}) \right\} / (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) , \quad (107)
\end{aligned}$$

with

$$c_{7ijk} = G_{1jG^+}^R a_{ji}^{\tilde{d}_k c} \quad c_{8ijk} = G_{1jG^+}^L a_{ji}^{\tilde{d}_k c} . \quad (108)$$

We have

$$\begin{aligned}
\Gamma_{\tilde{d} \tilde{d} W^+}(m_{\tilde{t}_1}^2) &= g^3 \sum_{j=1,2} \sum_{k=1\dots 3} \sum_{i=1,2} \frac{-G_{W^+ \tilde{t}_1 \tilde{d}_{ki}}}{\sqrt{2}} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* \left\{ [-a_{1i}^{d_k} (m_{\tilde{t}_1}^2 + m_{\tilde{\chi}_1^0}^2) - b_{1i}^{d_k} m_{\tilde{\chi}_1^0} m_{d_k}] \right. \\
&B_0(m_{\tilde{t}_1}^2, m_{\tilde{d}_{ki}}, M_W) + [2a_{1i}^{d_k} m_{\tilde{\chi}_1^0}^2 + b_{1i}^{d_k} m_{\tilde{\chi}_1^0} m_{d_k}] B_0(m_{\tilde{\chi}_1^0}^2, m_{d_k}, m_{\tilde{d}_{ki}}) \\
&+ 2a_{1i}^{d_k} (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) B_0(0, m_{d_k}, M_W) + [2a_{1i}^{d_k} (m_{\tilde{\chi}_1^0}^2 (m_{\tilde{\chi}_1^0}^2 - m_{\tilde{t}_1}^2 - m_{\tilde{d}_{ki}}^2) - \frac{1}{2} m_{d_k}^2 + M_W^2) \\
&+ m_{\tilde{t}_1}^2 (m_{\tilde{d}_{ki}}^2 - \frac{1}{2} m_{d_k}^2)) - b_{1i}^{d_k} m_{\tilde{\chi}_1^0} m_{d_k} (2m_{\tilde{t}_1}^2 + m_{d_k}^2 - M_W^2 - 2m_{\tilde{\chi}_1^0}^2)] \\
&\left. C_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_1^0}^2, 0, M_W, m_{\tilde{d}_{ki}}, m_{d_k}) \right\} / (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) . \quad (109)
\end{aligned}$$

And finally

$$\begin{aligned}
\Gamma_{\tilde{d} \tilde{\chi}^+ W^+} &= g^3 \sum_{j=1,2} \sum_{k=1\dots 3} \sum_{i=1,2} G_{W^+ \tilde{t}_1 \tilde{d}_{ki}} \mathcal{V}_{cd_k} \mathcal{V}_{td_k}^* \left\{ [c_{9ijk} m_{\tilde{t}_1}^2 + c_{10ijk} m_{\tilde{\chi}_1^0} m_{\tilde{\chi}_j^+}] B_0(m_{\tilde{t}_1}^2, M_W, m_{\tilde{d}_{ki}}) \right. \\
&+ [c_{9ijk} (m_{\tilde{\chi}_1^0}^2 - 2m_{\tilde{t}_1}^2) - c_{10ijk} m_{\tilde{\chi}_1^0} m_{\tilde{\chi}_j^+}] B_0(m_{\tilde{\chi}_1^0}^2, m_{\tilde{\chi}_j^+}, M_W) + [c_{9ijk} (m_{\tilde{t}_1}^2 m_{\tilde{\chi}_j^+}^2 - 2m_{\tilde{t}_1}^2 m_{\tilde{d}_{ki}}^2 \\
&+ m_{\tilde{\chi}_1^0}^2 m_{\tilde{d}_{ki}}^2) + c_{10ijk} m_{\tilde{\chi}_1^0} m_{\tilde{\chi}_j^+} (-m_{\tilde{t}_1}^2 + m_{\tilde{\chi}_j^+}^2 - m_{\tilde{d}_{ki}}^2 + m_{\tilde{\chi}_1^0}^2)] \\
&\left. C_0(m_{\tilde{t}_1}^2, m_{\tilde{\chi}_1^0}^2, 0, m_{\tilde{d}_{ki}}, M_W, m_{\tilde{\chi}_j^+}) \right\} / (m_{\tilde{t}_1}^2 - m_{\tilde{\chi}_1^0}^2) , \quad (110)
\end{aligned}$$

with

$$c_{9ijk} = G_{1jW^+}^R a_{ji}^{\tilde{d}_k c} \quad c_{10ijk} = G_{1jW^+}^L a_{ji}^{\tilde{d}_k c} . \quad (111)$$

## C FCNC counterterm

We start from the  $\tilde{u} - u$ -neutralino part of the Lagrangian in the interaction basis, expressed in terms of the bare squark and quark fields,  $\tilde{u}_i^{(0)}$  and  $u_i^{(0)}$  and the bare quark mass matrix  $\mathbf{m}_{ij}^{(0)}$ , where  $i, j = 1, 2, 3$  denote the generation indices and  $l = 1, \dots, 4$  the neutralino mass eigenstates,

$$\begin{aligned}
\mathcal{L}_{\tilde{u} \tilde{u} \tilde{\chi}^0} &= -\bar{u}_i^{(0)} g e_{Li}^{u_i} \tilde{u}_{iL}^{(0)} \mathcal{P}_R \tilde{\chi}_l^0 + \bar{u}_i^{(0)} \left( -\frac{g Z_{l4} \mathbf{m}_{ij}^{(0)}}{\sqrt{2} M_W \sin \beta} \right) \tilde{u}_{jR}^{(0)} \mathcal{P}_R \tilde{\chi}_l^0 \\
&\quad -\bar{u}_i^{(0)} g e_{Ri}^{u_i} \tilde{u}_{iR}^{(0)} \mathcal{P}_L \tilde{\chi}_l^0 + \bar{u}_i^{(0)} \left( -\frac{g Z_{l4} \mathbf{m}_{ij}^{(0)}}{\sqrt{2} M_W \sin \beta} \right) \tilde{u}_{jL}^{(0)} \mathcal{P}_L \tilde{\chi}_l^0 + h.c. . \quad (112)
\end{aligned}$$

The couplings  $e_{L,Rl}^{u_i}$  have been defined in Eq. (64). Let us look at the right-chiral part of the coupling. Rotation to the mass eigenstates yields

$$\begin{aligned} \mathcal{L}_{\bar{u}\tilde{u}\tilde{\chi}^0}^R &= -\bar{u}_k^{m(0)} U_{ki}^{uL(0)} g e_{Ll}^{u_i} \tilde{W}_{is}^{(0)\dagger} \tilde{u}_s^{m(0)} \mathcal{P}_R \tilde{\chi}_l^0 + \bar{u}_k^{m(0)} U_{ki}^{uL(0)} \frac{-g Z_{l4} \mathbf{m}_{ij}^{(0)}}{\sqrt{2} M_W \sin \beta} \tilde{W}_{j+3s}^{(0)\dagger} \tilde{u}_s^{m(0)} \mathcal{P}_R \tilde{\chi}_l^0 + h.c. \\ &\quad (i, j, k = 1, 2, 3, s = 1, \dots, 6). \end{aligned} \quad (113)$$

Note, that  $\tilde{W}_{is}^{(0)\dagger} \equiv \tilde{W}_{Lis}^{(0)\dagger}$ ,  $\tilde{W}_{j+3s}^{(0)\dagger} \equiv \tilde{W}_{Rjs}^{(0)\dagger}$ , cf. Eq. (10). Upon renormalization we replace [16]

$$\bar{u}^{m(0)} U^{uL(0)} \rightarrow \bar{u}^m \left( 1 + \frac{\delta Z^{L\dagger}}{2} \right) (1 + \delta u^{uL}) U^{uL} \quad (114)$$

$$\tilde{W}_{L,R}^{(0)\dagger} \tilde{u}^{m(0)} \rightarrow \tilde{W}_{L,R}^\dagger (1 + \delta \tilde{w}^\dagger) \left( 1 + \frac{\delta Z^{\tilde{u}}}{2} \right) \tilde{u}^m \quad (115)$$

$$\mathbf{m}^{(0)} \rightarrow \mathbf{m} + \delta \mathbf{m}, \quad (116)$$

where we have suppressed the indices. The wave function  $\tilde{u}^m$  denotes a six-component column vector. With the replacement  $\tilde{W}_{L,R} = W_{L,R} U^{uL,R}$ , cf. Eq. (15), we have for the Yukawa part of the coupling

$$\bar{u}^m \left( 1 + \frac{\delta Z^{L\dagger}}{2} \right) (1 + \delta u^{uL}) U^{uL} (\mathbf{m} + \delta \mathbf{m}) U^{uR\dagger} (1 + \delta u^{uR\dagger}) W_R^\dagger (1 + \delta \tilde{w}^\dagger) \left( 1 + \frac{\delta Z^{\tilde{u}}}{2} \right) \tilde{u}^m, \quad (117)$$

times  $(-g Z_{l4})/(\sqrt{2} M_W \sin \beta)$ . For the mass renormalization we choose the renormalization prescription such that the bare mass matrices and hence  $\delta \mathbf{m}$  are diagonal, *i.e.*

$$(1 + \delta u^{uL}) U^{uL} (\mathbf{m} + \delta \mathbf{m}) U^{uR\dagger} (1 + \delta u^{uR\dagger}) = (\mathbf{m}^D + \delta \mathbf{m}^D), \quad (118)$$

where  $D$  denotes diagonal matrices. This is possible since the off-diagonal elements can be absorbed into the off-diagonal elements of the antihermitian part of the right-handed wave function renormalization matrices [36]. Exploiting the unitarity of the mixing matrices we finally find for the renormalized Lagrangian in the mass eigenstate basis

$$\mathcal{L}_{\bar{u}\tilde{u}\tilde{\chi}^0} = \bar{u}_i^m (G_{isl}^R + \delta G_{isl}^R) \mathcal{P}_R \tilde{u}_s^m \tilde{\chi}_l^0 + \bar{u}_i^m (G_{isl}^L + \delta G_{isl}^L) \mathcal{P}_L \tilde{u}_s^m \tilde{\chi}_l^0 + h.c., \quad (119)$$

with the couplings given by

$$G_{isl}^R = -g e_{Ll}^{u_i} (W_L^\dagger)_{is} - \frac{g Z_{l4} m_{u_i} \delta_{ij}}{\sqrt{2} M_W \sin \beta} (W_R^\dagger)_{js} \quad (120)$$

$$G_{isl}^L = -g e_{Rl}^{u_i} (W_R^\dagger)_{is} - \frac{g Z_{l4} m_{u_i} \delta_{ij}}{\sqrt{2} M_W \sin \beta} (W_L^\dagger)_{js} \quad (121)$$

$$\begin{aligned} \delta G_{isl}^R &= -g e_{Ll}^{u_i} \left[ \frac{\delta Z_{ij}^{L\dagger}}{2} (W_L^\dagger)_{js} + (W_L^\dagger)_{it} \frac{\delta Z_{ts}^{\tilde{u}}}{2} + \delta u_{ij}^{uL} (W_L^\dagger)_{js} + (W_L^\dagger)_{it} \delta \tilde{w}_{ts}^\dagger \right] \\ &\quad - \frac{g Z_{l4}}{\sqrt{2} M_W \sin \beta} \left[ \frac{\delta Z_{ij}^{L\dagger}}{2} m_{u_j} \delta_{jk} (W_R^\dagger)_{ks} + m_{u_i} \delta_{ij} (W_R^\dagger)_{jt} \frac{\delta Z_{ts}^{\tilde{u}}}{2} \right. \\ &\quad \left. + m_{u_i} \delta_{ij} (W_R^\dagger)_{jt} \delta \tilde{w}_{ts}^\dagger + \delta m_{u_i} \delta_{ij} (W_R^\dagger)_{js} \right] \end{aligned} \quad (122)$$

$$\begin{aligned}
\delta G_{isl}^L = & -g e_{Rl}^{u_i} \left[ \frac{\delta Z_{ij}^{R\dagger}}{2} (W_R^\dagger)_{js} + (W_R^\dagger)_{it} \frac{\delta Z_{ts}^{\tilde{u}}}{2} + \delta u_{ij}^{u_R} (W_R^\dagger)_{js} + (W_R^\dagger)_{it} \delta \tilde{w}_{ts}^\dagger \right] \\
& - \frac{g Z_{l4}}{\sqrt{2} M_W \sin \beta} \left[ \frac{\delta Z_{ij}^{R\dagger}}{2} m_{u_j} \delta_{jk} (W_L^\dagger)_{ks} + m_{u_i} \delta_{ij} (W_L^\dagger)_{jt} \frac{\delta Z_{ts}^{\tilde{u}}}{2} \right. \\
& \left. + m_{u_i} \delta_{ij} (W_L^\dagger)_{jt} \delta \tilde{w}_{ts}^\dagger + \delta m_{u_i} \delta_{ij} (W_L^\dagger)_{js} \right] . \tag{123}
\end{aligned}$$

In the framework of MFV at  $\mu_{\text{MFV}}$  the  $W$  matrix is diagonal in flavour space at tree level. At one-loop flavour off-diagonal elements are induced through the wave function and the mixing matrix renormalization.

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